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UNIVERSITY OF ALBERTA

HYDRODYNAMICS OF UV LASER-PRODUCED PLASMA

BY

YING YIN TSUI



A thesis submitted to the Faculty of Graduate Studies and
Research in partial fulfillment of the requirements for the
degree of Doctor of Philosophy.

DEPARTMENT OF ELECTRICAL ENGINEERING

Edmonton, Alberta

Fall 1992



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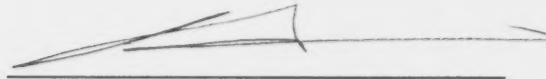
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ABSTRACT

Energy balance, plasma temperature, ablation parameters and asymptotic behavior of ions are measured for plasma produced by laser beams at wavelengths of 0.25 μm and 0.69 μm and pulse lengths of 1 ns to 25 ns with irradiances of $10^{10} - 10^{14} \text{ W/cm}^2$.

The results are consistent with inverse bremsstrahlung being the dominant absorption process. Energy balance and temperature intensity scaling are in reasonable agreement with analytic and numerical calculations. Scaling laws for the mass ablation rate and ablation pressure from aluminum plasma based on ion diagnostics are obtained at 0.25 μm for intensities ranging from $10^{10} - 10^{13} \text{ W cm}^{-2}$. The results are in reasonable agreement with a 1 D steady state ablation model and results from other laboratories. Energy transport measurement for gold plasma indicate the results are consistent with a two region plasma model consisting of a high temperature electron thermal transport region and a low temperature high density radiation transport region.

Detailed charge state distributions of the asymptotic plasma measured using a single shot electrodynamic analyzer and comparison with results from a hydrodynamic code that include time dependent ionization and recombination physics are reported for the first time. Reasonable agreement was found between the measured distributions and those calculated by recombination kinetics. However, discrepancies observed in detailed charge state spectrum and partial failure to predict the two peak behavior in aluminum may indicate inaccuracies in the recombination model and failure to include radiation energy transport in the modeling of the aluminum plasma.

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CHAPTER ONE

INTRODUCTION

The study of plasmas produced by submicron wavelength lasers is important both to obtain a better understanding of the basic processes taking place and to utilize them in a number of application areas. These applications include inertial confinement fusion (ICF) and soft x-ray sources to name a few. In the direct drive approach to ICF, fusion is initiated by the compression of a fuel pellet by laser ablation of its surface, whereas, in indirect x-ray driven fusion, x-rays are first produced in the high atomic number outershell of the pellet and the radiation is used to compress an inner fuel pellet. Both incoherent and coherent x-ray sources can be generated by laser produced plasmas. In x-ray lasers, plasmas are used as a gain medium. In recombination lasers the rapid expansion of these plasmas leave some atoms in highly excited states causing population inversion. In collisionally excited lasers electron collisions continuously pump the population inversion. These plasmas are also efficient emitter of ions and x-rays. Such ion and x-ray sources can be used in basic science, medical science and engineering.

When a laser beam is focused on matter it will ionize the atoms and produce a plasma which expands outwards. The laser energy is absorbed in the less dense part of the plasma known as the corona. The continuing absorption of laser energy will increase the temperature of the plasma creating a thermal gradient. Electron thermal conduction will carry heat into the material and hence create and ablate more plasma. The high temperature of the plasma implies that atoms are excited and ionized. As the plasma

expands outward its temperature decreases and the internal states of the ions change. Electromagnetic radiation is emitted from the plasma which partially is reabsorbed and partially escapes. As the plasma continues to expand most of its internal energy will be converted into kinetic energy. During the expansion the temperature and the density of the plasma drop and some of the highly ionized ions will fall to the next lower ionization state and release their ionization energy which causes the temperature to drop at a slower rate. However, the rapid decrease in density of the expanding plasma causes some atoms to be trapped in highly ionized states. These frozen charge states will relax to their equilibrium states at a very slow rate. Measurements made of the plasma at a large distance, i.e. centimeters, from the target surface will record the states of these ions after these closely related phenomena took place.

The goals of this thesis are to characterize the ion emission in detail, to determine the hydrodynamic energy and momentum balance in the laser ablation of targets and to investigate the ion recombination process.

In this thesis we report for the first time the detailed charge distributions of the asymptotic UV laser produced plasma measured using a single shot electrodynamic charge analyzer and comparison of experimental data to results from a hydrodynamic code that includes time dependent ionization and recombination physics.

In the experiments a wide range of experimental conditions are used to identify dominant mechanisms. Laser beams at wavelengths of $0.25\text{ }\mu\text{m}$ and $0.69\text{ }\mu\text{m}$ and pulse lengths of 1 ns to 20 ns with irradiances between 10^{10} - 10^{14} W/cm^2 are used in the experiments. The main target materials used in the experiments are carbon, aluminum and gold. The overall energy absorption is measured using an integrating sphere. The distribution of energy into ion kinetic energy and x-ray energy is measured by using plasma

calorimeters and x-ray calorimeters respectively. The temperature of the plasma is measured by metal foil filter techniques. The asymptotic behavior of the ion velocities and charge state spectrum is measured by using charge collectors and charge analyzers. Analytical and numerical models are used throughout this study to identify and quantify the important mechanisms.

This thesis is organized in eight chapters. Chapter one is the introduction. Chapter two reviews the theoretical models that are used throughout this thesis. Analytical models will be introduced first, hydrodynamic codes, both one dimensional and two dimensional will be introduced next and a one dimensional hydrodynamic code with recombination kinetics will be introduced last. Chapter three describes the laser systems and the main experimental apparatus used in this study. In chapter four the asymptotic behavior of the expanding plasmas will be reported. In chapter five the energy balance and plasma condition measurements are reported and analyzed. Chapter six presents measurements of ablation parameters and the energy transport measurements of the high atomic number material. Finally, discussion of results and conclusions of this thesis will be presented in chapters seven and eight.

CHAPTER TWO

THEORETICAL MODELS

While the physical world is not a continuum, plasmas produced by lasers can be usually described as a continuum due to the large number of particles that are present. In addition, the macroscopic view (continuum) and microscopic view (particles) are connected to each other, since properties in the macroscopic view can be interpreted as averages of events involving a large number of microscopic particles.

Because of the highly inter-related nature of the mechanisms involved in the interaction processes of lasers and matter, numerical simulations are essential in order to provide detailed solutions to the problems involved. However, analytical solutions are useful to isolate dominant physical mechanisms and to provide simple scaling laws. They can give a better understanding of the physics involved and provide a means to understand the results of more complicated simulations.

In this chapter, the fluid equations that commonly are used to describe the hydrodynamics of the plasmas produced by lasers will be reviewed. Brief descriptions of hydrodynamics codes Castor and Medusa will follow. Two simple models which describe the recombination kinetics will be described. Finally, some analytic models will be presented.

2.1 Fluid Equations

Three conservation laws describe the physics of plasmas produced by lasers. The continuity equation describes the conservation of material density,

$$\frac{\partial \rho}{\partial t} + \vec{\nabla} \cdot (\rho \vec{v}) = 0 \quad (2.1)$$

where ρ and $\vec{v}$ are the density and velocity of the plasma. The momentum equation describes the conservation of momentum,

$$\frac{\partial \rho \vec{v}}{\partial t} + \vec{\nabla} \cdot (\rho \vec{v} \vec{v}) + \vec{\nabla} p = 0 \quad (2.2)$$

where p is the pressure of the plasma. The energy equation describes the conservation of energy,

$$\frac{\partial E}{\partial t} + \vec{\nabla} \cdot (\vec{v} E) + \vec{\nabla} \cdot \vec{q} + \vec{\nabla} \cdot (p \vec{v}) = W \quad (2.3)$$

$$E = \rho(e + \frac{1}{2} v^2) \quad (2.4)$$

where e and $\vec{q}$ are internal energy and heat flux of the plasma and W is the energy change per unit volume due to an external source such as the absorbed laser energy.

The electrons and ions of the plasma are usually described separately and connected via terms that describe the coupling of ions and electrons. However, for sufficiently high temperature and/or low densities (as in the corona of the laser-produced plasmas) where the heating mechanisms, the energy transport and their effects are not very dependent on the ion parameters, a single fluid set of equations using the mass and momentum of the cold ions and pressure and temperature of the hot electrons can be used.

2.2 Equations of State

To solve the fluid equations (2.1) to (2.3) a knowledge of the equation of state $e = e(p, \rho)$ is required. Hydrodynamic codes which perform simulations of laser-plasma interactions require the use of equation of

state models which describe the thermodynamic functions of electrons and ions. Different models exist for different regime of densities and temperatures. In general, a hydrodynamic code will use a suitable model in the regimes of relevance and join the different regimes by semi-empirical interpolation methods. In the corona the plasma can be described by ideal gas laws

$$e = \frac{3}{2} p \quad (2.5)$$

$$p = (n_e + n_i) k_B T_e \quad (2.6)$$

where n_e , n_i , k_B , T_e are the number densities for electrons and ions, Boltzmann constant and electron temperature respectively.

2.2 Three Dimensional Simulations

Ideally three dimensional (3D) simulations should be used to model the physics of the plasmas produced by lasers. However, detailed 3D simulations are rare since they require an enormous amount of computer power to run. With the rapid advancement of computer technology 3D simulations of fluid dynamics are starting to be actively pursued.

2.3 Two Dimensional Simulations

Two dimensional (2D) simulations can be used to model phenomena that show symmetry. Plasmas produced by lasers focused onto planar targets can be approximated as being cylindrically symmetric about an axis perpendicular to the target surface. A 2D hydrodynamics code Castor^{1,2} is used to simulate experiments in this study. The simulation results are provided by Dr. R. Rankin and Dr. C. Capjack at

the University of Alberta. The Castor calculations use cylindrical Eulerian co-ordinates to calculate the behavior of laser-irradiated targets of various atomic numbers. This code does not include detailed atomic physics. Thus it cannot correctly describe phenomena dominated by radiation transport and recombination kinetics.

2.4 One Dimensional Simulations

Strictly speaking one dimensional (1D) codes are limited in their ability to reproduce experiments that focus laser beams onto a planar target. Compared to 2D simulations, 1D simulations are simpler in nature and the results are easier to interpret. Thus, the results can be useful in understanding those limiting cases where the plasma expansion may be considered as 1D. In addition, detailed atomic physics and ionization and recombination kinetics algorithms can be included in a 1D hydrodynamics code to shed light on the physics of these important processes.

2.4.a Medusa

Medusa³ is a 1D hydrodynamics code that uses Lagrangian co-ordinates. A radiation transport algorithm⁴ has been included in this code. This code can be run with or without the radiation transport turned on. By comparing the two results the effect of radiation transport on various processes can be estimated.

2.4.b. Simple Hydrocode with Recombination Kinetics

A 1D hydrodynamic code including recombination physics that uses Eulerian co-ordinates has been developed as part of the thesis work to help explain experimental observations of expanding plasma at a large distance from

the target surface. This code solves the three hydrodynamic equations in cartesian, cylindrical, and spherical geometry as well as the ionization and recombination kinetic equations. The hydrodynamic equations in cartesian, cylindrical, and spherical geometry where $\alpha = 1, 2, 3$ respectively are

$$\frac{\partial \rho}{\partial t} + \frac{1}{R^{\alpha-1}} \frac{\partial}{\partial R} (R^{\alpha-1} \rho v) = 0 \quad (2.7)$$

$$\frac{\partial \rho v}{\partial t} + \frac{1}{R^{\alpha-1}} \frac{\partial}{\partial R} (R^{\alpha-1} \rho v^2) + \frac{\partial p}{\partial R} = 0 \quad (2.8)$$

$$\frac{\partial E}{\partial t} + \frac{1}{R^{\alpha-1}} \frac{\partial}{\partial R} [R^{\alpha-1} (v E + v p)] = W \quad (2.9)$$

Assume that the ions are distributed in a number of ionization levels and let f_i be the fraction of ions that populate a level i , where level i has a charge of $(i-1)$ then the density can be expressed as

$$\rho = \sum_i \rho_i = \sum_i f_i \rho \quad (2.10)$$

The total derivative of the density following the fluid particles for each level can be written as

$$\frac{D \rho_i}{D t} = \frac{\partial \rho_i}{\partial t} + \frac{1}{R^{\alpha-1}} \frac{\partial}{\partial R} (R^{\alpha-1} \rho_i v) \quad (2.11)$$

using equations (2.7) and (2.10) it follows

$$\sum_i \frac{D \rho_i}{D t} = \sum_i \rho \frac{D f_i}{D t} = 0 \quad (2.12)$$

where the total derivative of ionization fraction following a fluid particle is given by

$$\frac{D f_i}{D t} = \frac{\partial f_i}{\partial t} + v \frac{\partial f_i}{\partial R} \quad (2.13)$$

As shown in figure 2.1 the population of a level i can be increased or decreased via its own or its neighboring levels' ionization and recombination reactions. An equation that describes the evolution of the charge states in the local fluid particle frame of reference is given by

$$\frac{D f_i}{D t} = n_e (S_{i-1} f_{i-1} - (S_i + \alpha_i) f_i + \alpha_{i+1} f_{i+1}) \quad (2.14)$$

where S_i and α_i are the ionization and recombination rates of level i respectively and n_e is the electron number density. As the plasma expands its temperature will drop and some of its ions will recombine to a lower level and thus release heat back to the plasma. Let E_{ret} be the heat returned to the plasma, then equation (2.9) becomes

$$\frac{\partial E}{\partial t} + \frac{1}{R^{\alpha-1}} \frac{\partial}{\partial R} [R^{\alpha-1} (v E + v p)] = E_{ret} \quad (2.15)$$

The ideal gas law is used as the equation of state of the plasma. In summary, the equations used in the code are the following:

$$\frac{\partial \rho}{\partial t} + \frac{1}{R^{\alpha-1}} \frac{\partial}{\partial R} (R^{\alpha-1} \rho v) = 0 \quad (2.7)$$

$$\frac{\partial \rho v}{\partial t} + \frac{1}{R^{\alpha-1}} \frac{\partial}{\partial R} (R^{\alpha-1} \rho v^2) + \frac{\partial p}{\partial R} = 0 \quad (2.8)$$

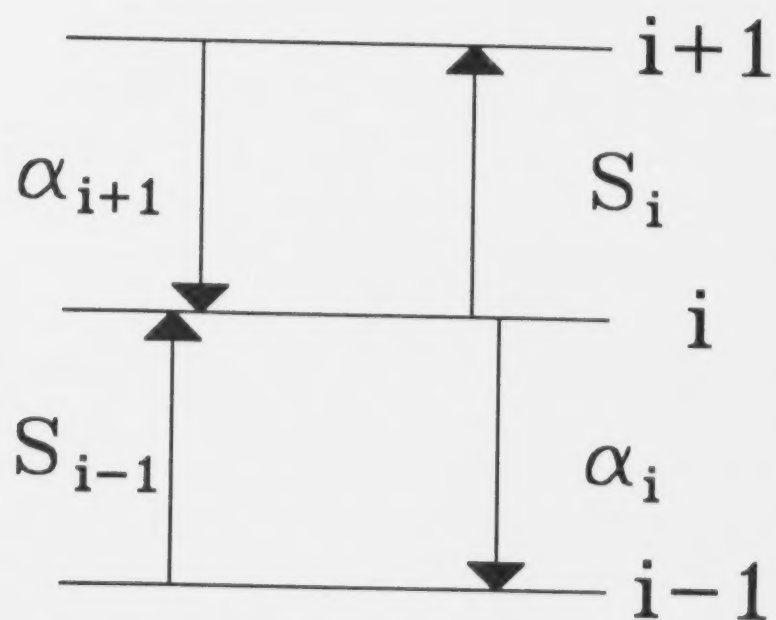


Figure 2.1 Schematic diagram to illustrate how the population of a charge state level i can be changed due to the ionization and recombination transitions from state i and its neighbor states $i-1$ and $i+1$.

$$\frac{\partial E}{\partial t} + \frac{1}{R^{\alpha-1}} \frac{\partial}{\partial R} [R^{\alpha-1} (v E + v p)] = E_{ret} \quad (2.15)$$

$$E = \frac{3}{2} p + \frac{1}{2} \rho v^2 \quad (2.16)$$

$$\frac{\partial f_i}{\partial t} + v \frac{\partial f_i}{\partial R} = n_e (S_{i-1} f_{i-1} - (S_i + \alpha_i) f_i + \alpha_{i+1} f_{i+1}) \quad (2.14)$$

The calculation of plasma flow produced by heating a planar target with a finite focal spot is a two dimensional problem. The plasma diverges due to the finite size of the focal spot. If the laser pulse is long enough (≥ 1 ns) a quasi-stationary flow will be established. The flow is approximately spherically diverging over most of the area of the spot with a center of curvature one focal spot diameter behind the target surface⁵ as shown in figure 2.2. Thus a 1D spherical code can be used as an approximation to the real problem.

2.4.b.1 Ionization and Recombination Rate Coefficients

The electron impact ionization rate from the ion ground state, averaged over a Maxwellian electron distribution, can be expressed as⁶

$$S_i = \frac{10^{-5} \zeta_i \left(\frac{T_e}{E_\infty^i} \right)^{\frac{1}{2}}}{(E_\infty^i)^{\frac{3}{2}} \left(6 + \frac{T_e}{E_\infty^i} \right)} \exp \left(\frac{-E_\infty^i}{T_e} \right) \quad [\text{cm}^3 \text{s}^{-1}] \quad (2.17)$$

where E_∞^i , T_e and ζ_i are the ionization energy from state i to $i+1$, electron temperature in eV and number of outer

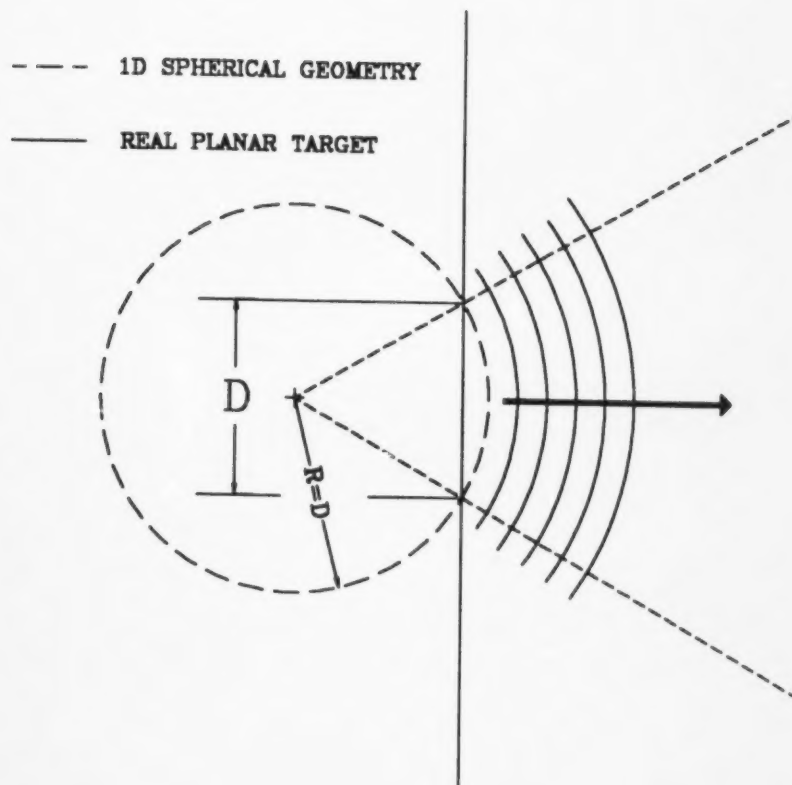


Figure 2.2 The planar target is modeled by an equivalent spherical target. Here D is the diameter of the laser focal spot.

shell electrons respectively. The recombination rate α_i consists of the radiative recombination rate α_{Ri} and three body recombination α_{3Bi} rate through the relation

$$\alpha_i = \alpha_{Ri} + n_e \alpha_{3Bi} \quad (2.18)$$

For the radiative recombination one has⁷

$$\begin{aligned} \alpha_{R(1+1)} &= 5.2 \times 10^{-14} i \left(\frac{E_\infty^i}{T_e} \right)^{\frac{1}{2}} \left\{ 0.43 + 0.5 \ln \left(\frac{E_\infty^i}{T_e} \right) + 0.47 \left(\frac{E_\infty^i}{T_e} \right)^{-\frac{1}{3}} \right\} \\ &\quad [\text{cm}^3 \text{s}^{-1}] \quad (2.19) \end{aligned}$$

For the three body recombination coefficient one has^{8,9,10}

$$\begin{aligned} \alpha_{3B(1+1)} &= 8.75 \times 10^{-27} i^3 T_e^{-\frac{9}{2}} [\ln(i^2 + 1)]^{\frac{1}{2}} \\ &\quad [\text{cm}^6 \text{s}^{-1}] \quad (2.20) \end{aligned}$$

The above expressions for the recombination coefficients are valid only for hydrogen-like atoms when the electron temperature is much lower than the ionization potential.

2.4.b.2 Energy Returned to the Plasma

In the recombination of an electron with an ion in a three body collision, the electron is first captured by the ion into one of the upper levels of the atom with a binding energy $\Delta E_{\infty, n} \approx kT_e$. In general, as shown in figure 2.3 the distance between the levels $\Delta E_{n, n-1}$ is much less than kT_e . For these dense levels, collision transitions to the neighboring levels are the most probable. An electron in an energy level j in this regime has almost equal probability to undergo a transition to levels $j+1$ or $j-1$. Thus the bound electron moves diffusively down the energy levels of

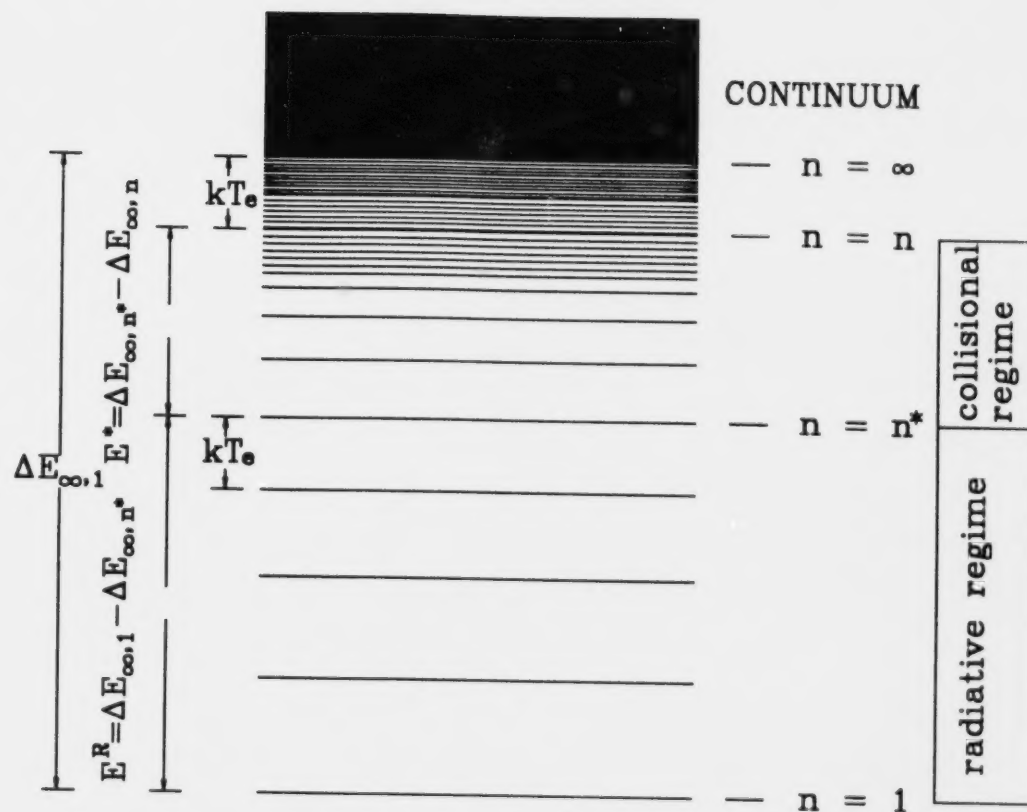


Figure 2.3 Schematic diagram to illustrate the downward transition of an electron capture by an atom during a three body recombination event.

the atom towards the ground level. When the electron reaches a level $n = n^*$ where $\Delta E_{n^*, n^*-1} = kT_e$, spontaneous radiative transitions take over as the more probable mechanism. The energy $E^* = \Delta E_{\infty, n^*} - \Delta E_{\infty, n}$ released in the de-excitation process by the collision transitions returns to the plasma as heat. The fraction of the line radiation energy $E^R = \Delta E_{\infty, 1} - \Delta E_{\infty, n^*}$ returned to the plasma depends on the opacity of the plasma. The corona of the plasma, because of its low density, is transparent to radiation. As a good approximation, one can assume that the energy E^R , along with energy released by radiative recombination, leaves the plasma. Thus, the quantity E^* represents the mean energy per recombination event which is transferred to the free electrons. To evaluate E^* one can estimate and compare the rates of collision and radiation transition of the electron through the atomic energy levels. Kuznetsov and Raiser⁹ have derived E^* for a hydrogenic atom

$$E_{(1+1)}^* = 1.1 \times 10^{-3} \bar{Z}^{\frac{1}{6}} i^{\frac{2}{3}} n_i^{\frac{1}{6}} T_e^{\frac{1}{12}} \quad [\text{eV}] \quad (2.21)$$

This then implies the heat return to the plasma can be expressed¹¹ as

$$E_{\text{ret}} = \sum_i n_e^2 n_i \alpha_{3Bl} f_i E_i^* \quad [\text{eV cm}^{-3}] \quad (2.22)$$

2.4.b.3 Boundary and Initial Conditions

The recombination code uses initial hydrodynamic conditions from the 2D hydro code, Castor, and then calculates the evolution of ion states over the long path to the experimental ion detectors. The left hand boundary is chosen at $R = R_L = 184 \mu\text{m}$, approximately two focal spot diameters away from the target surface and the right hand boundary at $R = R_R = 1\text{m}$. Low density gas ($5 \times 10^{15} \text{ cm}^{-3}$) in

a box is assumed as the initial condition ($t = 0$) of the code to avoid numerical instability. The code uses the constant density, velocity and temperature initial conditions from the Castor simulation for the region from $R = 184 \mu\text{m}$ to $640 \mu\text{m}$ with an exponential decay for $R > 640 \mu\text{m}$

$$\rho(R, t=0) = \begin{cases} \rho_0 & \text{for } R \leq R_L \\ \rho_0 \exp\left(-\frac{(R-R_L)}{R_L}\right) & \text{for } R > R_L \end{cases} \quad (2.23)$$

$$v(R, t=0) = 0 \quad (2.24)$$

$$T(R, t=0) = \begin{cases} T_0 & \text{for } R \leq R_L \\ T_0 \exp\left(-\frac{(R-R_L)}{R_L}\right) & \text{for } R > R_L \end{cases} \quad (2.25)$$

The initial charge distribution is assumed to be neutral

$$f_1 = 1 \quad (2.26)$$

$$f_i = 0 \quad \text{for } i \neq 1 \quad (2.27)$$

For $0 < t < 6 \text{ ns}$ the quantities ρ , v , T from the 2D code Castor at $R = 184 \mu\text{m}$ are read in as time dependent boundary conditions at the left hand boundary. Since Castor does not include detailed atomic physics, no charge state distribution is available. The time dependent boundary conditions for f_1 are obtained by post processing the density and temperature data from Castor to solve the recombination and ionization rate kinetic equations in Lagrangian co-ordinates from target surface to the left hand boundary.

For $6 < t < 40 \text{ ns}$ the quantities ρ , v , T from the 2D code Castor at $R = 184 \mu\text{m}$ are linearly extrapolated from the output at $t = 5.5 \text{ ns}$ and 6 ns and used as time dependent boundary conditions at the left hand boundary. The left hand boundary condition for f_1 is assumed to be

constant in time given by the output of f_i at $t = 6$ ns from integrating the ionization and recombination rate kinetic equations.

For $t > 40$ ns, the left hand boundary conditions for ρ , v , T and f_i are equated to 10^{-10} to avoid divisions by zero.

For the right hand boundary a free boundary condition is assumed

$$\frac{\partial \rho}{\partial R} = 0 \quad (2.28)$$

$$\frac{\partial \rho v}{\partial R} = 0 \quad (2.29)$$

$$\frac{\partial E}{\partial R} = 0 \quad (2.30)$$

$$\frac{\partial f_i}{\partial R} = 0 \quad i=1, n \quad (2.31)$$

2.4.b.4 Numerical methods

The system of non-linear partial differential equations including equations (2.3), (2.4), (2.12), (2.13) and (2.14) are solved by the numerical package MOL1D¹². MOL1D developed by J. M. Hyman at Los Alamos Laboratory is a FORTRAN subroutine package for the method of lines solution for systems of initial-boundary-value partial differential equations in one spatial dimension.

2.5 Simple Recombination Kinetics Model for Uniform Spherical Plasma

A simpler model has also been developed to describe spherical freely expanding plasmas. In laser produced plasmas the hydrodynamic variables are complicated

functions of time and space. Consider an expanding laser produced plasma which was created in a short interval of time sufficiently long ago that the front boundary expands freely into vacuum at constant velocity. This can be described by a model consisting of many spherical shell elements each having its own density and velocity. Each of these elements can be described as a gas occupying initially, at $t = t_0$, a spherical volume of radius R_0 with uniform density ρ_0 expanding into vacuum at a constant velocity. This situation will occur when most of the plasma's internal energy has been converted into kinetic energy and the plasma flow becomes collisionless^{10,11}. The flow velocity of the plasma will approach an asymptotic value, that is, the plasma reaches a state of free expansion. The equations describing a freely expanding plasma element are the recombination kinetic equation

$$\frac{d f_1}{d t} = n_e (S_{1-1} f_{1-1} - (S_1 + \alpha_1) f_1 + \alpha_{1+1} f_{1+1}) \quad (2.32)$$

and the energy equation modified by the heat returned to the plasma

$$\frac{d e}{d t} + p \frac{d n_1^{-1}}{d t} = E_{ret} \quad (2.33)$$

Since the velocity of the element is time independent, the radii of the boundaries of the element must be proportional to time. From this it follows that the volume of the element is proportional to the cubic power of time giving

$$n_1(t) = n_1(t=t_0) \left(\frac{t_0}{t} \right)^3 \quad (2.34)$$

The set of equations can be completed by using the ideal

gas law equations

$$e = \frac{3}{2} p \quad (2.35)$$

$$p = (n_e + n_i) k T_e \quad (2.36)$$

The above system of nonlinear ordinary differential equations are solved by the numerical package GEARB¹³. GEARB, developed by A. C. Hindmarsh at the Lawrence Livermore Laboratory, is a FORTRAN subroutine package for the solution of the initial value problem for systems of ordinary differential equations with a Jacobian matrix which has banded or nearly banded form.

2.6 Analytic Models

It is possible to construct simple analytic models of the interaction of a laser with a slab target. Such models are useful as they help isolate the dominant mechanisms and give useful simple scaling laws. They also provide a means to understand the results from more complicated computer simulations.

The flow geometry of plasmas produced by lasers is 2D. However, it can be described initially as a 1D planar self-similar flow followed by a 1D spherical stationary flow. This description of the plasma flow is supported by interferometric measurement of electron density of these plasmas¹⁴ as well as results from Castor simulations. Depending on the laser focal spot size and laser pulse length used, the spherical or the planar flow can be the dominant flow pattern.

It is useful to compare results from 1D analytical models with those from 1D numerical simulations. Most analytic models are only good for approximate estimations within a factor of about two. By comparing the two results

the accuracy of a particular analytic model can be determined.

2.5.a Isothermal Plane Self-Similar Expansion

A simple analytic solution¹¹ describing the expanding plasma can be obtained from the continuity equation and momentum conservation equation if the following assumptions are used

- I. planar geometry
- II. isothermal plasma
- III. cold ions
- IV. ideal gas equation of state

Because the plasma is isothermal and the ions are cold, the pressure gradient term $\partial P / \partial x$ in equation (2.2) can be re-written as $C_s^2 \partial \rho / \partial x$. The acoustic speed of the plasma is given by

$$C_s = \left(\frac{\bar{Z}}{A m_p} k_B T_e \right)^{\frac{1}{2}} \quad (2.37)$$

Where A and m_p are the mass number and proton mass respectively and $\bar{Z}$ is the average charge. The two fluid equations (2.1) and (2.2) become two coupled partial differential equations in variables ρ and V . These two equations describe self-similar motion of the plasma in the sense that they contain no characteristic time or length. Thus a transformation with $\xi = x/t$ as a new variable can be used to transform the two partial differential equations into two ordinary differential equations. The solutions to these equations for $x > -C_s t$ are

$$V = C_s + \frac{x}{t} \quad (2.38)$$

$$\rho = \tilde{\rho} \exp - \left(1 + \frac{x}{C_s t} \right) \quad (2.39)$$

where $\tilde{\rho}$ is the unperturbed plasma density. The solutions describe a rarefaction wave propagating into the undisturbed hot plasma at the acoustic speed of the plasma (see figure 2.4).

The isothermal assumption is not unreasonable, since high temperature in the plasma corona implies that the thermal conductivity is high. The self-similar solution is valid for the region of expanding plasma where its Debye length is insignificant compared to any other length in the process. The ions are assumed to be cold due to expansion cooling. The ion temperature is not important as long as the number of electrons greatly exceeds the number of ions. If the ions are assumed to have the same temperature as the electrons the only modification to the solution is a slight change in the acoustic speed; the acoustic speed would differ by a factor of $\sqrt{(Z+1)/Z}$.

Equations (2.32) and (2.33) give a density profile extending to infinity. As mentioned above these solutions are not valid for the region of the expanding plasma where the Debye length is comparable to a physical length in the process. This implies that the plasma will be ended by a region of charge separation¹⁵.

2.5.b Isothermal Spherical Stationary Expansion

A simple analytic solution^{14,16} describing the expanding plasma can be obtained from the continuity equation and momentum conservation equation if the following assumptions are used

- I. spherical geometry
- II. isothermal plasma
- II. stationary supersonic flow
- IV. velocity of the flow is C_s at $r = r_0$

The solutions for this problem can be derived by setting the time derivatives of the continuity and momentum equations (2.7) and (2.8) to zero and the velocity equal to

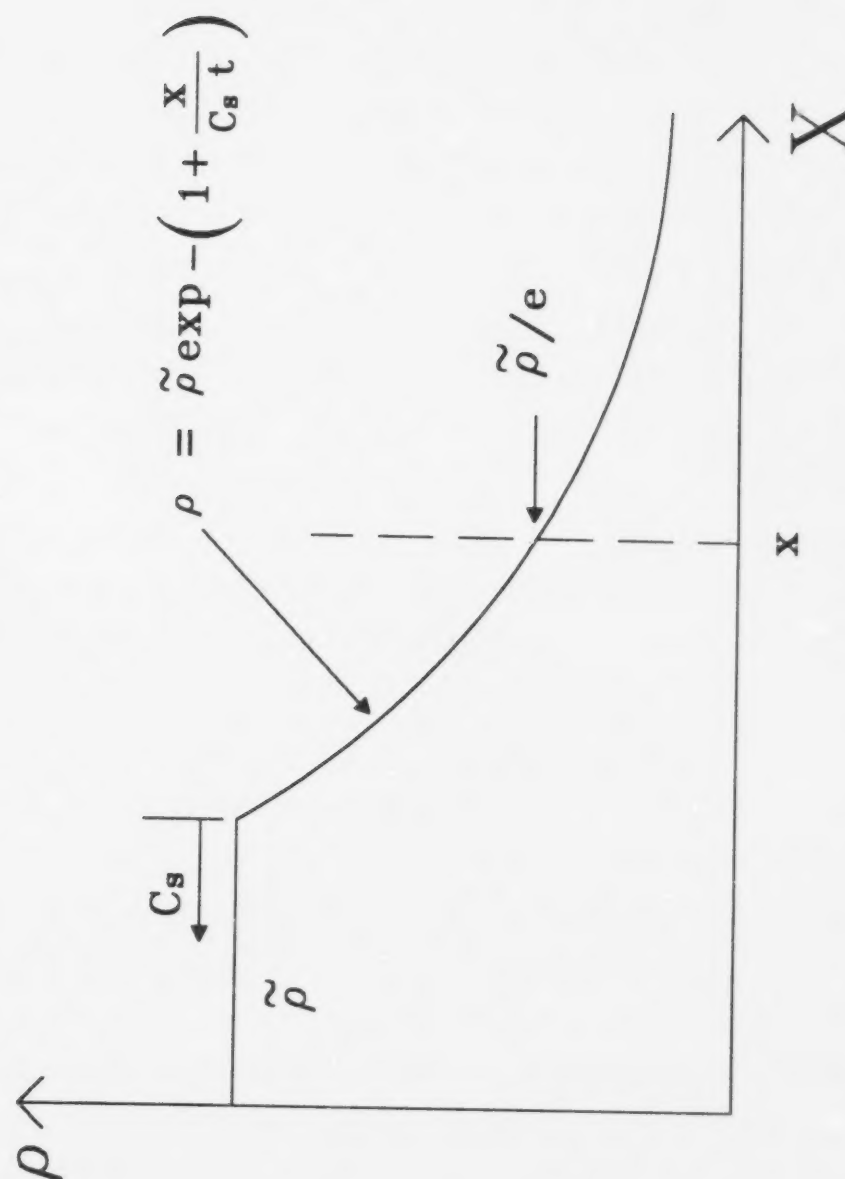


Figure 2.4 Planar isothermal self-similar expansion into vacuum.

the acoustic velocity at $r = r_0$. The asymptotic result for the region $r \gg r_0$ is given by

$$\rho = \frac{\rho_0 C_s}{V} \left(\frac{r_0}{r} \right)^2 \quad (2.40)$$

$$V = C_s \left(1 + 4 \ln \left(\frac{r_0}{r} \right) \right)^{\frac{1}{2}} \quad (2.41)$$

2.5.c Steady State Ablation Model

A steady state ablation model¹⁷ can be derived with some simple assumptions. There are four distinct regions in this model (see figure 2.5):

Region I

The less dense plasma corona where the laser energy is absorbed up to the absorption surface.

Region II

The conduction zone bounded by the absorption surface and the ablation surface. The absorbed laser energy penetrates this region solely by heat conduction.

Region III

The solid compressed and heated by the shock created by the reaction force to the plasma ablation.

Region IV

The unperturbed solid material ahead of the shock front.

By using some simple yet reasonable hypotheses scaling laws for temperature (T_e), mass ablation rate ($\dot{m}$), ablation

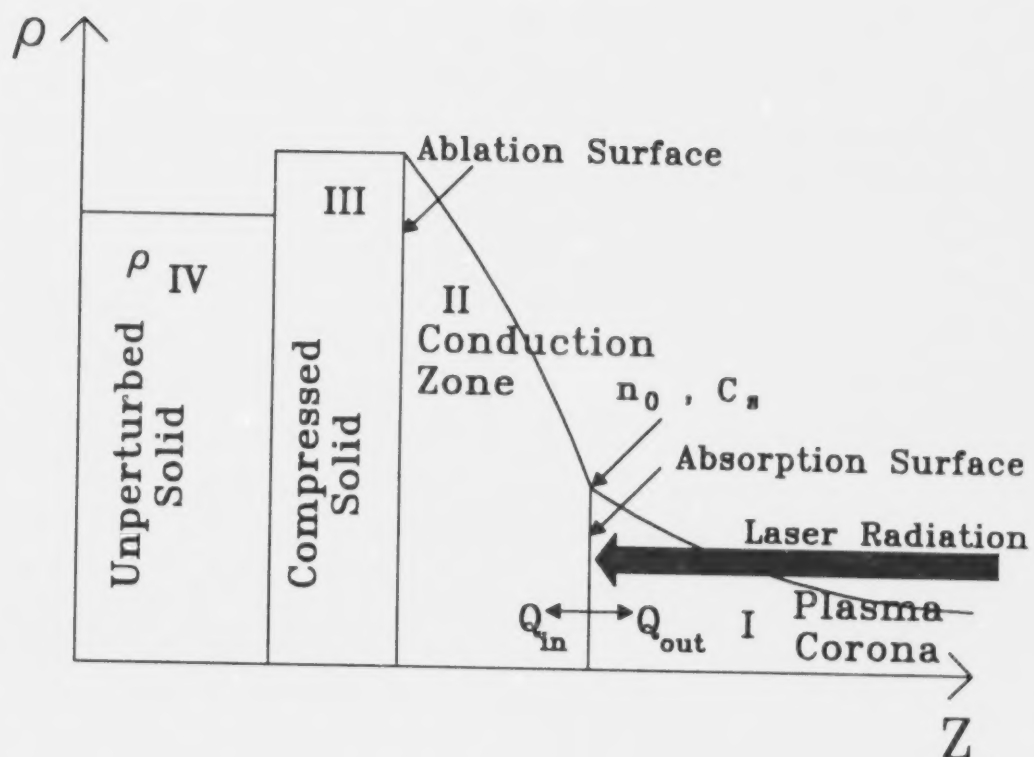


Figure 2.5 Self-regulating plasma flow model

pressure (P_a), and penetration density (n_0) can be derived. These hypotheses are listed below:

- I. The laser intensity is constant during the duration of the laser pulse τ . The laser energy is absorbed by inverse bremsstrahlung. At the characteristic penetration density n_0 the plasma flows at the ion acoustic velocity C_s and the laser radiation has attenuated to $\exp(-2)$ of its original value.
- II. The plasma expansion in the corona can be described by an isothermal plane self-similar expansion model or an isothermal spherical stationary expansion model.
- III. The equation of state of the plasma is assumed to be the same as an ideal gas.
- IV. The conduction zone has reached a steady state and solutions for the plasma flow in the corona can be joined smoothly to the solutions for the plasma flow in this region.

In the inverse bremsstrahlung process laser light is absorbed by causing oscillation of the electrons that collide with ions thereby converting coherent energy of oscillation into random thermal energy. The absorption of laser energy by inverse bremsstrahlung through a plasma is given by¹⁷

$$A = 1 - \exp\left(-\int_{x_0}^{\infty} a_{\omega}(x) dx\right) \quad (2.42)$$

where

$$a_{\omega}(x) = \eta \bar{Z} \ln \Lambda T_e^{-\frac{3}{2}} n_c \left[\frac{n(x)}{n_c}\right]^2 \left[1 - \frac{n(x)}{n_c}\right]^{-\frac{1}{2}} \quad (2.43)$$

$\bar{Z}$ is the average charge state of the plasma and the coefficient in front is given by

$$\eta = \frac{e^4}{3 (2 \pi)^{\frac{3}{2}} (m_e c^2)^{\frac{1}{2}} \epsilon_0^2 k_B^{\frac{3}{2}}} \quad (2.44)$$

The Coulomb logarithm for our regime of interest is given by

$$\ln \Lambda = \ln \left(\frac{3 (4 \pi \epsilon_0)^{\frac{3}{2}}}{2 \bar{Z} e^3} \left(\frac{k_B^3 T_e^3}{\pi n_e} \right)^{\frac{1}{2}} \right) \quad (2.45)$$

and the critical density of the plasma is given by

$$n_c = \frac{4 \pi^2 m_e \epsilon_0 c^2}{e^2 \lambda_L^2} \quad (2.46)$$

The maximum density, n_0 , that the laser beam can effectively penetrate to can be derived from equation (2.42) if the density profile of the plasma in the corona is known. For both the planar and stationary spherical flow cases for densities well below critical the density n_0 can be derived using hypothesis II together with equations (2.38) and (2.39), assuming a $1/r^2$ density profile for the spherical case.

Planar Case

$$n_0 = \frac{2 n_c^{\frac{1}{2}} T_e^{\frac{3}{4}}}{\eta^{\frac{1}{2}} (\bar{Z} \ln \Lambda \tau)^{\frac{1}{2}} C_s^{\frac{1}{2}}} \quad (2.47)$$

Spherical Case

$$n_0 = \frac{6^{\frac{1}{2}} n_c^{\frac{1}{2}} T_e^{\frac{3}{4}}}{\eta^{\frac{1}{2}} (\bar{Z} \ln \Lambda r_0)^{\frac{1}{2}}} \quad (2.48)$$

Since the model assumes that the corona is isothermal, it is not important where the laser light is absorbed exactly. For convenience, the laser light will be assumed to be fully absorbed at the density n_0 . The internal energy density and pressure of the plasma is given by

$$e = \frac{3}{2} n_e k_B T_e = \frac{3}{2} \rho C_s^2 \quad (2.49)$$

$$p = n_e k_B T_e = \rho C_s^2 \quad (2.50)$$

Thus the energy conservation equation for the plasma is

Planar Case

$$\begin{aligned} \frac{\partial}{\partial t} \left(\frac{3}{2} \rho C_s^2 + \frac{1}{2} \rho v^2 \right) + \frac{\partial}{\partial x} \left[\rho v \left(\frac{1}{2} v^2 + \frac{3}{2} C_s^2 \right) + v p + Q \right] \\ - I_A \delta(x - x_0) = 0 \end{aligned} \quad (2.51)$$

Spherical Case

$$\begin{aligned} \frac{\partial}{\partial t} \left(\frac{3}{2} \rho C_s^2 + \frac{1}{2} \rho v^2 \right) \\ + \frac{1}{r^2} \frac{\partial}{\partial r} \left(r^2 \left[\rho v \left(\frac{1}{2} v^2 + \frac{3}{2} C_s^2 \right) + v p + Q \right] \right) \\ - I_A \delta(x - x_0) = 0 \end{aligned} \quad (2.52)$$

The heat flux (Q_{out}) flowing into the corona is determined in the planar case by integrating the energy conservation equation from the absorption surface to ∞ . The Q_{out} for the planar flow case is $\rho_0 C_s^3$. Assuming a stationary flow in the spherical case would require no outward heat flux. However, since the spherical flow pattern expands outwards into vacuum, energy must be added to the coronal plasma and it is assumed that the outward heat flux is the same as for the planar case¹⁷. The solution for the plasma flow in the corona region is thus determined.

The assumption that the conduction zone is in a steady

state is not unreasonable. A large class of plasmas produced by lasers can be described as self-regulating. In these plasmas the thickness of the conduction zone is approximately constant. If the thickness of the zone changes then the energy available to ablate the solid also changes, resulting in a change of the rate of plasma production. The result is a feedback mechanism that cause the region to reach a steady state. The solution for the plasma flow in the conduction zone can be easily found by using hypothesis IV in the governing conservation equations with time derivatives equated to zero,

Planar Case

$$\frac{\partial}{\partial x}(\rho v) = 0 \quad (2.53)$$

$$\frac{\partial}{\partial x}(p + \rho v^2) = 0 \quad (2.54)$$

$$\frac{\partial}{\partial x}[\rho v (\frac{1}{2} v^2 + \frac{3}{2} c_s^2) + v p + Q_{in}] - I_A \delta(x-x_0) = 0 \quad (2.55)$$

Spherical Case

$$\frac{1}{r^2} \frac{\partial}{\partial r}[r^2 (\rho v)] = 0 \quad (2.56)$$

$$\frac{1}{r^2} \frac{\partial}{\partial r}[r^2 (p + \rho v^2)] = 0 \quad (2.57)$$

$$\frac{1}{r^2} \frac{\partial}{\partial r}(r^2 [\rho v (\frac{1}{2} v^2 + \frac{3}{2} c_s^2) + v p + Q_{in}]) - I_A \delta(x - x_0) = 0 \quad (2.58)$$

The mass ablation rate ($\dot{m}$), ablation pressure (P_A) and the heat flux (Q_{in}) flowing into the conduction zone are obtained by integrating these equations, using hypothesis

IV assuming the velocity of the ablation is much less than C_s . They are the same for both planar and spherical cases, assuming the thickness of the conduction zone is thin compared to the radius for the spherical case.

Planar and Spherical Cases

$$\dot{m} = \rho_0 C_s \quad (2.59)$$

$$P_A = 2 \rho_0 C_s^2 \quad (2.60)$$

$$Q_{in} = 3 \rho_0 C_s^3 \quad (2.61)$$

Energy balance requires the absorbed laser intensity is equal to the total heat flux in the plasma. Thus

Planar and Spherical Case

$$I_A = Q_{in} + Q_{out} = 4 \rho_0 C_s^3 = 4 \frac{A}{Z} m_p n_0 C_s^3 \quad (2.62)$$

The scaling laws for the ablation parameters follow from equations (2.47, 2.48, 2.49, 2.59, 2.60, 2.62)

Planar Case

$$T_{ev} = 279.5 (I_{13} \lambda_{\mu m})^{\frac{1}{2}} \left(\frac{A}{Z} \right)^{\frac{1}{8}} (\bar{Z} \ln A \tau_{ns})^{\frac{1}{4}} \quad (2.63)$$

$$\dot{m} \text{ g cm}^{-2} \text{ s}^{-1} = 9.75 \times 10^4 I_{13}^{\frac{1}{2}} \lambda_{\mu m}^{-\frac{1}{2}} \left(\frac{A}{Z} \right)^{\frac{7}{8}} (\bar{Z} \ln A \tau_{ns})^{-\frac{1}{4}} \quad (2.64)$$

$$P_A \text{ Mbar} = 3.03 I_{13}^{\frac{3}{4}} \lambda_{\mu m}^{-\frac{1}{4}} \left(\frac{A}{Z} \right)^{\frac{7}{16}} (\bar{Z} \ln A \tau_{ns})^{-\frac{1}{8}} \quad (2.65)$$

$$n_0 \text{ cm}^{-3} = 3.42 \times 10^{21} I_{13}^{\frac{1}{4}} \lambda_{\mu m}^{-\frac{3}{4}} \left(\frac{A}{Z} \right)^{\frac{5}{16}} (\bar{Z} \ln A \tau_{ns})^{-\frac{3}{8}} \quad (2.66)$$

Spherical Case

$$T_{\text{ev}} = 80.1 (I_{13} \lambda_{\mu\text{m}})^{\frac{4}{9}} \left(\frac{A}{Z} \right)^{\frac{2}{9}} (\bar{Z} \ln \Lambda r_0 \mu\text{m})^{\frac{2}{9}} \quad (2.67)$$

$$\dot{m} \text{ g cm}^{-2} \text{ s}^{-1} = 3.26 \times 10^5 I_{13}^{\frac{5}{9}} \lambda_{\mu\text{m}}^{-\frac{4}{9}} \left(\frac{A}{Z} \right)^{\frac{7}{9}} (\bar{Z} \ln \Lambda r_0 \mu\text{m})^{-\frac{2}{9}} \quad (2.68)$$

$$P_{\Lambda} \text{ Mbar} = 5.71 I_{13}^{\frac{7}{9}} \lambda_{\mu\text{m}}^{-\frac{2}{9}} \left(\frac{A}{Z} \right)^{\frac{7}{18}} (\bar{Z} \ln \Lambda r_0 \mu\text{m})^{-\frac{1}{9}} \quad (2.69)$$

$$n_0 \text{ cm}^{-3} = 2.22 \times 10^{22} I_{13}^{\frac{1}{3}} \lambda_{\mu\text{m}}^{-\frac{2}{3}} \left(\frac{A}{Z} \right)^{\frac{1}{6}} (\bar{Z} \ln \Lambda r_0 \mu\text{m})^{-\frac{1}{3}} \quad (2.70)$$

where the quantities are given in the units indicated and I_{13} is the intensity in units of 10^{13} W/cm^2 .

Chapter Three

Laser Systems and Experimental Details

The Krypton Fluoride (KrF) laser, with a wavelength of 248 nm, is an ultraviolet (UV) excimer laser which appears to offer a number of properties which make it a good candidate as a driver of laser fusion: (1) it is inherently a high power laser, (2) it operates at a short wavelength which is expected to be nearly optimum for laser-plasma coupling, (3) its lower laser level is unstable and dissociates implying a high overall efficiency can be attained, and (4) it is a gas laser implying that a high repetition rate can be attained.

The main laser system used in this thesis work is the Krypton Fluoride (KrF) laser system developed in the Department of Electrical Engineering at the University of Alberta. This system was developed for laser fusion related studies. With pulse compression, this laser system is capable of producing laser pulses of approximately 2 ns duration (FWHM), which is the width of the laser pulse between its 50% peak intensity points. To extend the parameter space so wavelength and pulse length scaling could be studied two secondary laser systems were also used. The first was the KrF laser system run without the pulse compression system to produce laser pulses of approximately 20 ns duration (FWHM). The second was a solid state ruby laser system with a wavelength of 694 nm. This system produces pulses of approximately 20 ns duration (FWHM).

The experimental instruments used in this thesis work can be classified into three major areas: (1) those used for measuring energy balance of the plasma, (2) those used for measuring the temperature of the plasma and (3) those used for measuring the asymptotic behavior of the ablating plasma. An integrating sphere, x-ray calorimeter and plasma

calorimeters were used to study the energy balance of the plasmas. The metal foil filter x-ray detector technique was used to study the temperature of the plasmas. Charge collectors and charge analyzers were used to study the asymptotic behavior of the expanding plasma.

3.1 Short Pulse KrF Laser System

The KrF laser system consists of four major components. These are the discharge laser front end, the electron beam pumped power amplifier (e-beam amplifier), the optical pulse compressor and the target chamber.

A schematic diagram of the original discharge laser front end^{18,19} is shown in figure 3.1. It consists of a master oscillator which was equipped with a multi-etalon system to control the line width of the laser beam and three amplifiers A1, A2, A3. The system generates 1.5 joule (J) of narrow band (~3 GHz) radiation in pulses of 20 ns duration (FWHM). Part of the laser beam is used to trigger and extract energy from the e-beam amplifier. The remaining part of the laser beam pumps the optical pulse compressor and is used for pulse shortening by a nonlinear optical process.

The e-beam amplifier²⁰ is capable of generating 30 J of radiation in a pulse of 60 ns duration (FWHM). It supplies essentially all of the optical energy that is available for irradiation of the target. A schematic diagram of an e-beam amplifier is shown in figure 3.2. Two five stage Marx banks are electrically triggered to pulse charge water-filled Blumlein pulse forming lines. These subsequently create a 60 ns 450 kV pulse at the two electron guns upon closure of laser-triggered spark gap switches. The resulting electron beams pump the laser cell from opposite sides to give maximum uniformity of energy deposition. The vacuum diode electron guns have thin aluminum-coated-Kapton as anodes. The anode foil and a thin Inconel pressure foil are mounted

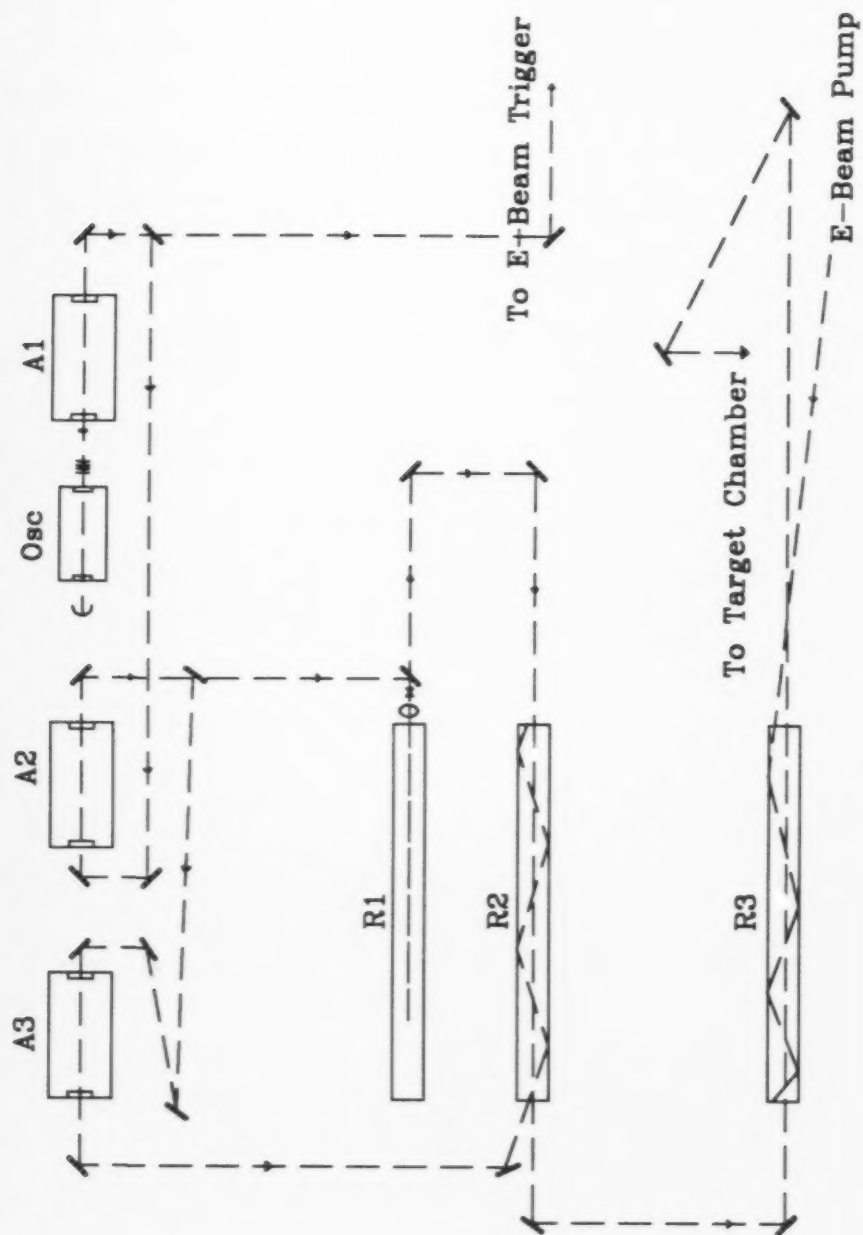


Figure 3.1 Schematic diagram of the KrF front end and Raman pulse compression systems.

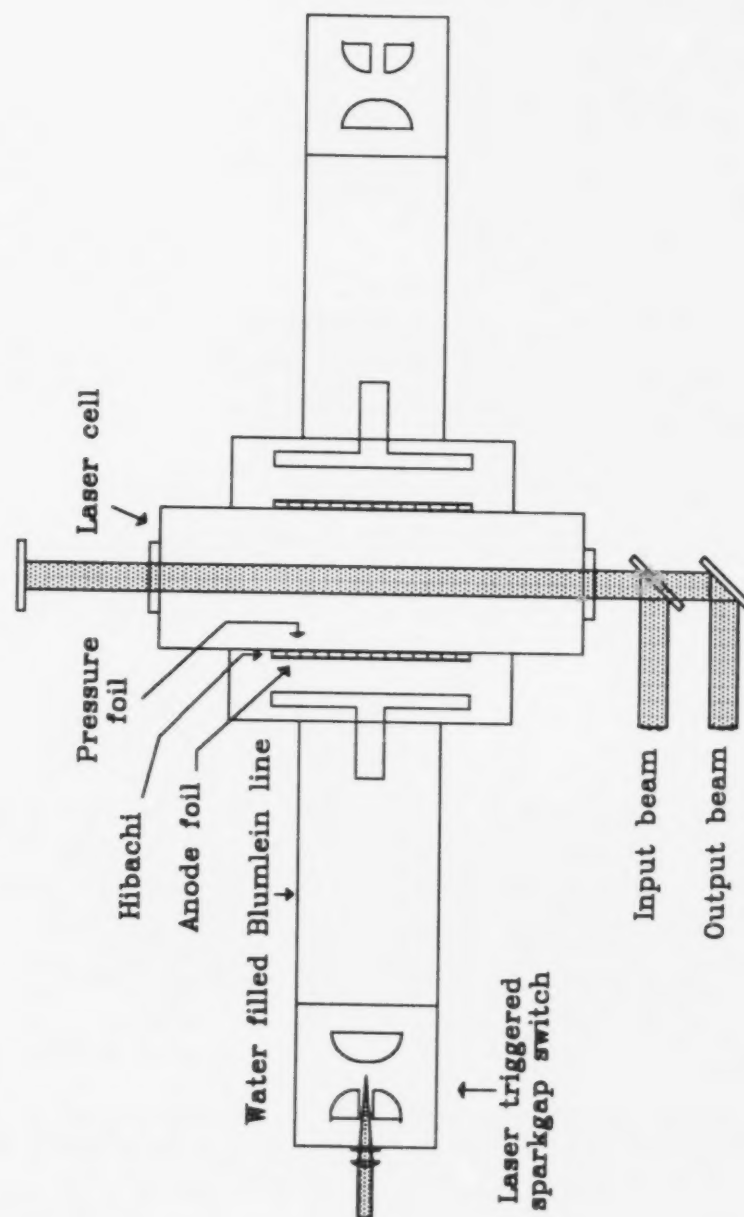


Figure 3.2 Schematic diagram of the e-beam amplifier.

on opposite sides of a steel grid (hibachi) and separate the diode from the high pressure laser gas mix since the diode requires vacuum ($<10^{-4}$ torr) for correct operation.

The KrF laser is a non-storage laser because of its short upper state lifetime and high stimulated emission cross-section. Instead, energy is continuously extracted for the duration of the excimer excitation. To generate the very short pulses that are required for target experiments, special techniques must be employed. Two different pulse compression systems, namely the Raman pulse compression system and Brillouin pulse compression system are used to produce the short pulses.

In the past, the Raman pulse compression system was used. It was operated using the backward stimulated Raman scattering process (SRS)²¹. The principle on which backward Raman pulse compression operates is shown in figure 3.3. In figure 3.3.b a long laser pump pulse propagates through a gaseous medium containing methane (CH_4). The pump photons with frequency ν_p can interact with a vibrational transition in the methane molecule with a frequency $\nu_m = 2917 \text{ cm}^{-1}$ generating a first Stokes frequency $\nu_s = \nu_p - \nu_m$ (figure 3.3.a). Just as the pump pulse reaches the exit window, a shorter pulse at frequency ν_s is injected into the cell in the opposite direction. As the two pulses move through one another, the Stokes pulse grows at the expense of the pump pulse. The net result is that a fraction of the energy from the long pulse is compressed into the short pulse ν_s . Figure 3.1 illustrates the set up of the Raman pulse compressor. It consists of three high pressure methane cells: R1, the Raman generation cell, R2, the preamplifier stage and R3, the power amplifier cell. The initial Stokes pulse is generated using backward SRS of part of the KrF laser pulse from A2 discharge module. The pump pulse for R1 is derived from A2 and the pump pulse for R2 is derived from A3. The pump pulse for the power

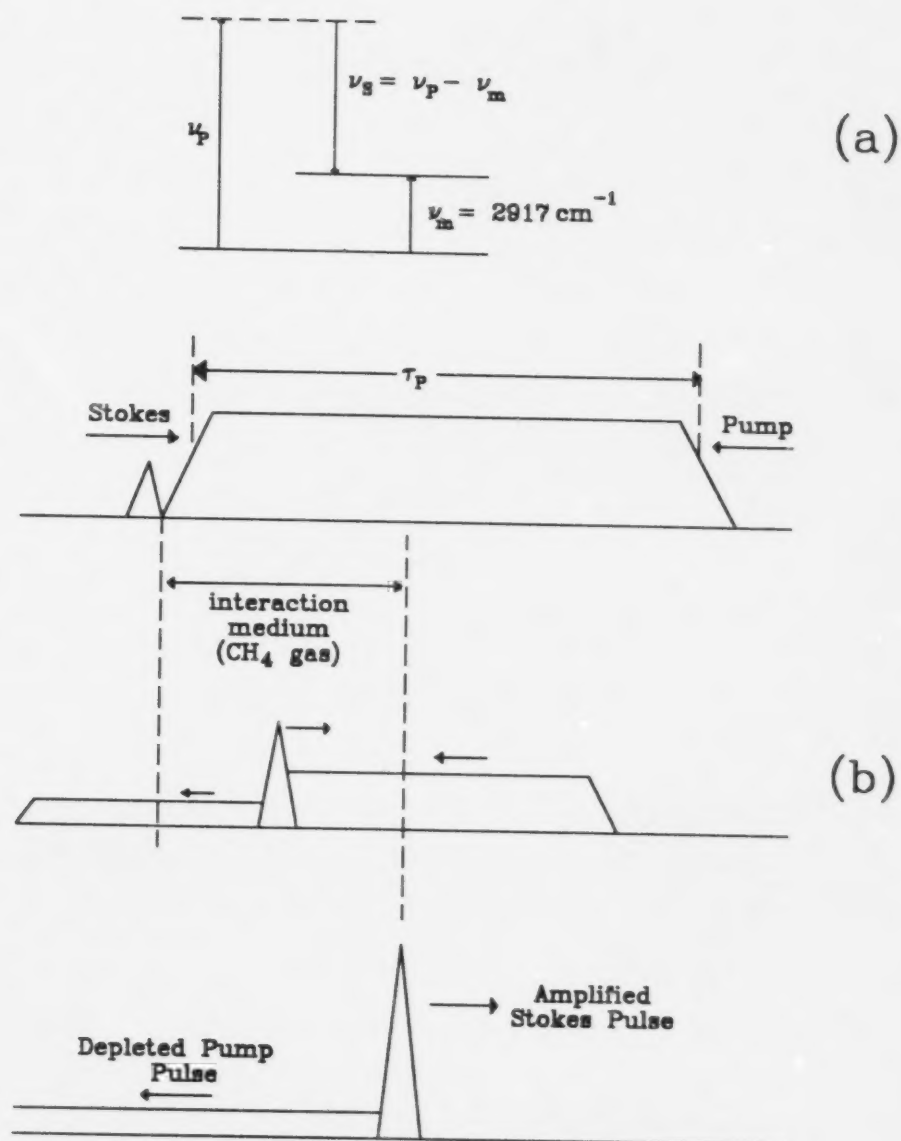


Figure 3.3 Backward Raman pulse compression.

amplifier stage, R3, is derived from the e-beam amplifier. The resulting laser pulse available on target was 2.0 J of 2.0 ns duration (FWHM) at a wavelength of 268 nm .

There are limits to the degree of pulse compression at high efficiency using stimulated Raman scattering because of build-up of second Stokes radiation²². The growth of the second Stokes pulse propagating together with the high intensity amplified first Stokes pulse can deplete the energy of the first Stokes pulse and hence reduce the energy conversion from the pump beam.

The present pulse compression system is based on an alternative technique, stimulated Brillouin Scattering (SBS). The principle on which backward Brillouin pulse compression operates is similar to backward Raman pulse compression. The major difference is that in SBS the scattering process is due to acoustic phonons rather than optical photons as in SRS. This implies the frequency shift $\nu_p - \nu_s$ in SBS is much less than in SRS. The Brillouin pulse compression system runs most efficiently²² when a pump laser band width of $\leq 100\text{MHz}$ is used. Such narrow line width is achieved by injection locking the KrF laser front end using the frequency doubled output from an argon ion/dye laser amplifier system pumped by a xenon chloride (XeCl) laser²³.

Figure 3.4 illustrate the Brillouin pulse compressor system. It consists of four high pressure SF_6 cells: B1, the Brillouin generation cell, B2, the preamplifier stage and B3 and B4, the power amplifier cells. The initial Stokes pulse is generated using backward SBS of part of the pulse from the A2 discharge module. The pump pulse for B1 is derived from A2 and the pump pulse for B2 is derived from A3. The pump pulses for the power amplifier stages, B3 and B4, are derived from the e-beam amplifier. The resulting laser pulse energy available on target was 2.0 J of 1.5 ns duration (FWHM) at a wavelength of 248 nm.

After pulse compression by SRS or SBS, the laser beam

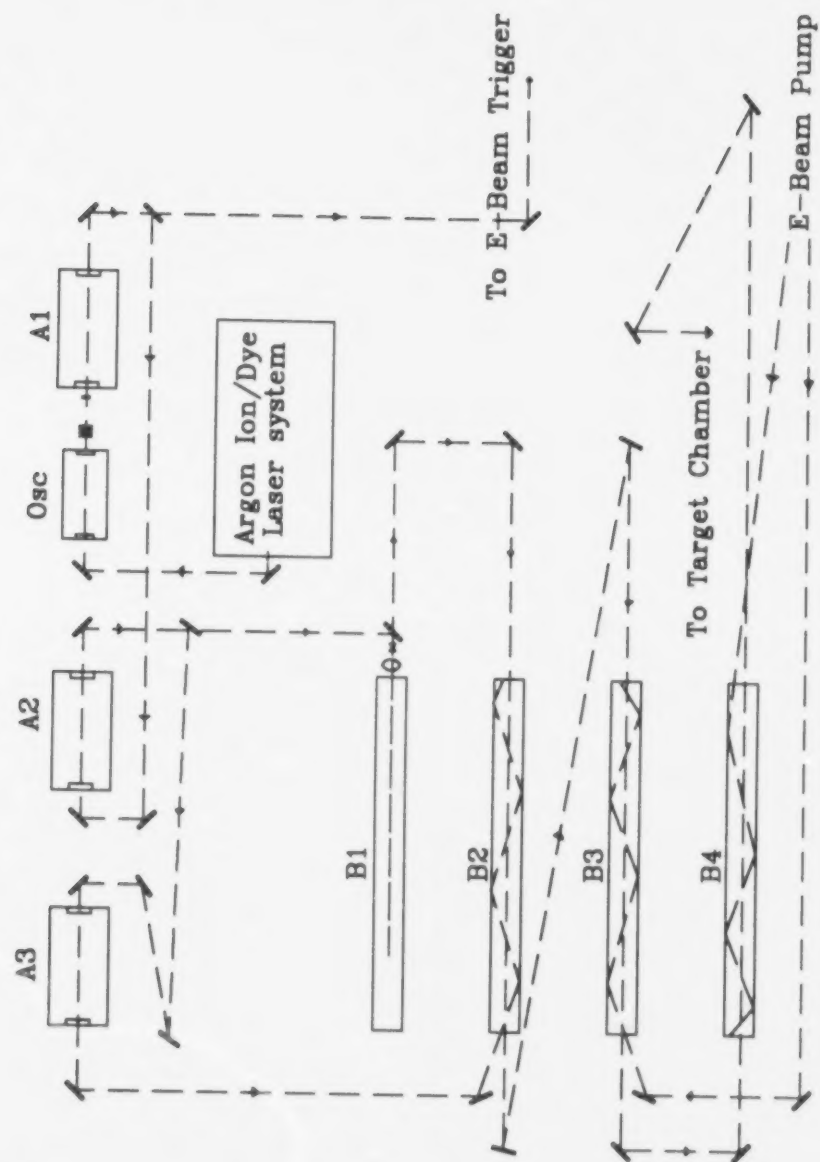


Figure 3.4 Schematic diagram of the KrF front end and Brillouin pulse compression systems.

is directed to a vacuum target chamber for the interaction experiments. Before entering the target chamber the beam parameters are diagnosed by different monitors as shown in figure 3.5. The temporal pulse shape is monitored by a Hamamatsu R1193U photodiode, PD, as well as a Hamamatsu C979 UV streak camera system. The laser energy on target is measured by a pyroelectric calorimeter CAL. The equivalent focal plane intensity is measured using a quartz lens, L2, which brings part of the beam to a focus. The focal spot distribution is subsequently imaged via lens L3 onto a UV vidicon camera camera and analyzed by means of a video digitizer system. The 4 cm x 6 cm laser beam is focused into the centre of the target chamber by means of a 10 cm diameter plano-convex 70 cm focal length quartz lens L1.

The target chamber which houses our laser-plasma interaction diagnostics is an octagonal aluminum/steel chamber measuring 80 cm across and 50 cm in height. It is equipped with a refrigerated alcohol cold trap and with a diffusion pump capable of producing a vacuum of nominally 10^{-6} torr inside the chamber. The target is located at the centre of the chamber using a motorized three axis positioning stage with external position read out and a position accuracy of 1 μm . Correct positioning of the target at the chamber centre is determined by using an imaging lens, L4, and observation screen, S, illuminated by a cw He-Ne alignment laser.

3.2 Long Pulse KrF Laser System

The output of A3, a 1.5 cm x 3 cm laser beam, is sent directly to the target chamber to provide a 400 mJ, 20 ns duration (FWHM) laser pulse at 248 nm.

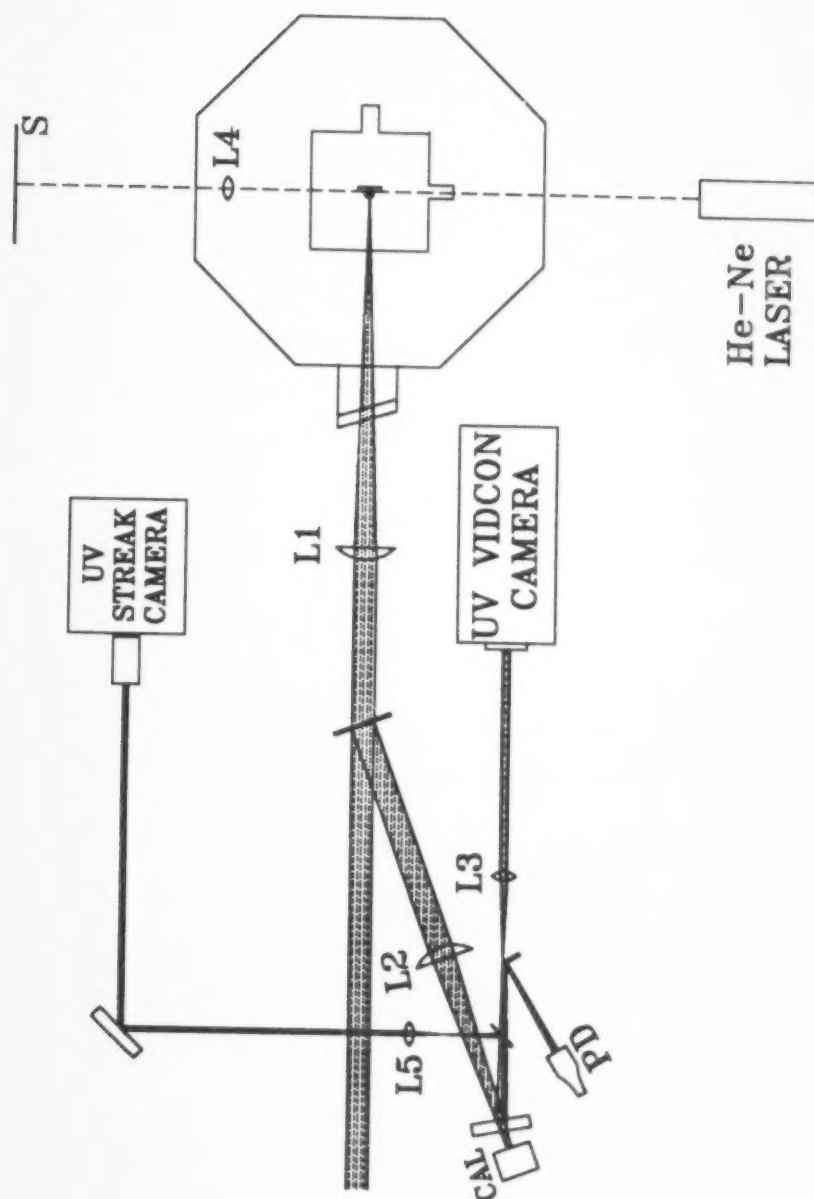


Figure 3.5 Schematic diagram of the beam diagnostics and target chamber of the KrF laser system.

3.3 Ruby Laser System

The solid state ruby laser system consists of three major components. These are the laser oscillator, laser amplifier and target chamber. A schematic diagram of the laser system is shown in figure 3.6. The laser oscillator consists of a 12.5 cm long by 8 mm diameter, A-R coated ruby laser rod that is pumped by two EG&G FX-47C-6.5 linear flash lamps inside a double elliptical polished silver coated reflector cavity. The flash lamp and laser rod are both water cooled to 15° C while the cavity is air cooled. The laser resonator is bounded by two mirrors M1 and M2 that are both flat with R=100% and R=40% dielectric coating respectively. A 6 mm-diameter iris limits the beam size. A Lasermetrics 1057FV pockels cell Q-switch and a calcite polarizer of Glan-Thompson type are used to produce the Q-switch pulse by half wave retardation. The laser amplifier consists of a 20 cm by 1.5cm diameter, ruby laser rod that is pumped by two EG&G FX-47C-6.5 linear flash lamps inside a double elliptical polished silver coated reflector cavity. The flash lamps and laser rod are both water cooled to 15° C while the cavity is air cooled. Both ends of the laser rod are cut at Brewster's angle to minimize loss of laser light due to reflection. The laser amplifier flashlamp current is delayed by 320 μ s relative to the laser oscillator flashlamp current to optimize amplification. The resulting laser pulse available on target has an energy of 1.5 J at a wavelength of 694 nm and a duration (FWHM) of 20 ns.

The laser beam is directed to a vacuum target chamber. Before entering the target chamber the beam conditions are diagnosed by different detectors as shown in figure 3.7. The temporal pulse shape is monitoring by a Hamamatsu R1193U photodiode, PD. The laser energy on target is measured by a pyroelectric calorimeter CA. The equivalent focal plane intensity is measured using a glass lens , L2,

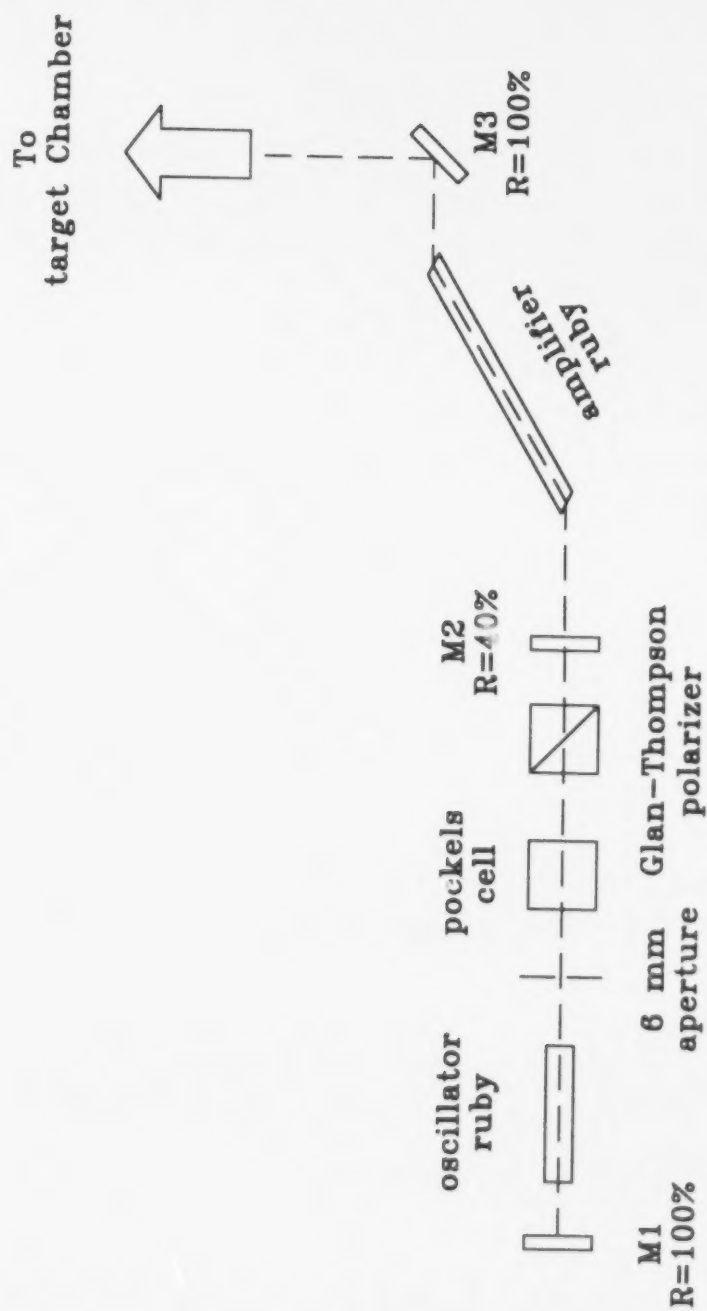


Figure 3.6 Schematic diagram of the ruby laser system.

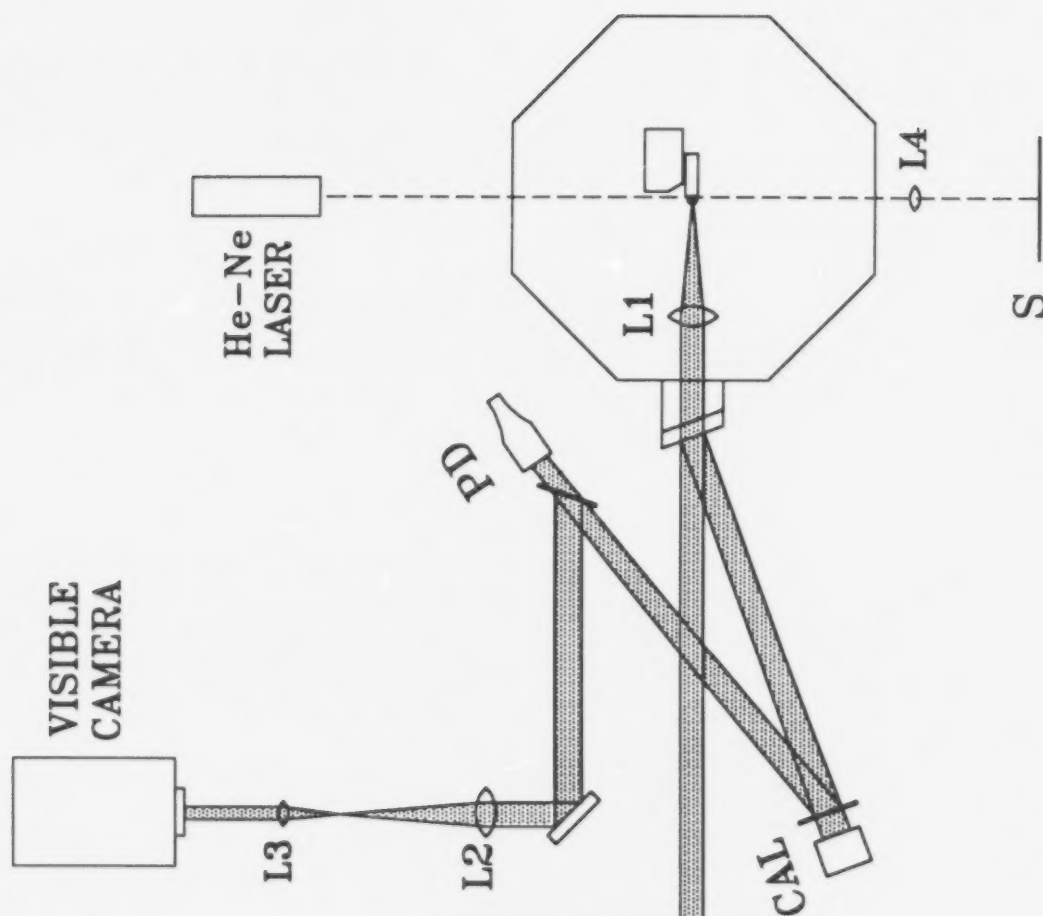


Figure 3.7 Schematic diagram of the beam diagnostics and target chamber of the ruby laser system.

which brings part of the beam to a focus. The focal distribution is subsequently imaged via lens L3 onto a visible vidicon camera. The 1.5 cm diameter circular laser beam is focused into the centre of the target chamber by means of a 5 cm bi-convex 14 cm focal length glass lens L1.

The target chamber which houses our laser-plasma interaction diagnostics is a polygonal aluminum chamber measuring about 50 cm across and 25 cm in height. It is equipped with a liquid nitrogen cold trap and a diffusion pump capable of producing a vacuum of $\approx 10^{-6}$ torr inside the chamber. The target is located at the centre of the chamber using a manually controlled three axis positioning stage. Targets are 1 cm thick disks with the curved surface facing the incoming laser beam. New target surface are obtained by rotation of the target. Correct positioning of the target at the chamber centre is determined by using an imaging lens, L4, and observation screen, S, illuminated by a cw He-Ne alignment laser.

3.4 Determination of Laser Focal Plane

The best focal condition is obtained with the following procedures. The position of the focusing lens is adjusted to within a mm of best focus condition of the laser by examining the laser burn pattern on a piece of UV sensitive Dylux instant image proof paper under a magnifier. The final adjustment of the position of the focusing lens to within the focal depth of the focus of the laser is determined as the point with the highest KeV energy x-ray production.

Characterization of a laser-plasma interaction experiment requires the measurement of the spatial and temporal distribution of the laser energy at the target. Typical time integrated spatial distributions at laser focus and the temporal distributions from the laser beams

from the short pulse KrF laser system with Brillouin and Raman compression schemes, the long pulse KrF laser system and the ruby laser system are shown in figure 3.8, 3.9, 3.10 and figure 3.11 respectively.

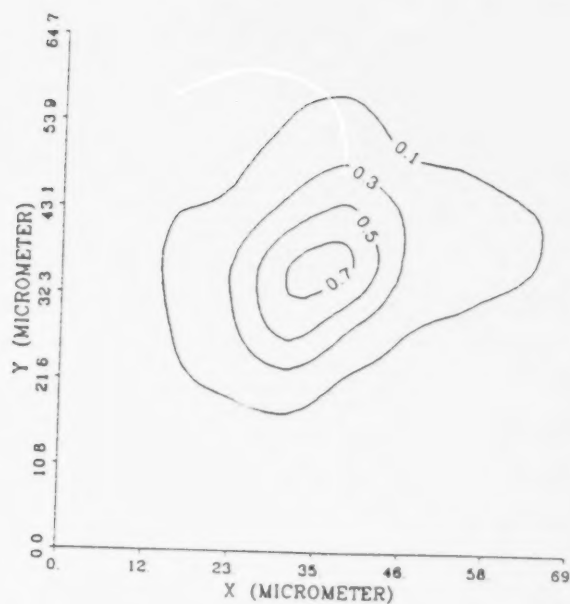
3.5 Summary of Laser Parameters

The KrF laser focusing systems are also run with the focus lens 2 mm (f1) and 4 mm (f2) away from its optimum position to increase the range of laser intensity available on target. Figure 3.12 gives the equivalent focal diameter of 90 % energy contour at various positions of the focusing lens. The effective laser intensity on target is calculated according to the following formula

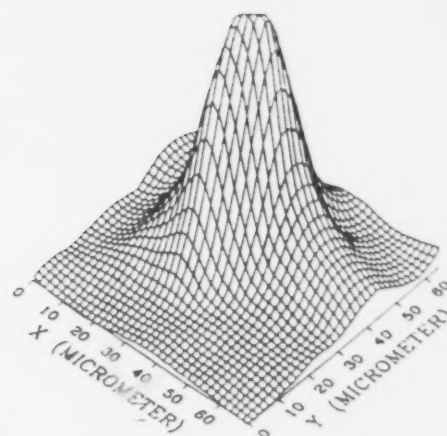
$$I_{nx} = \frac{E_{nx}}{\pi \left(\frac{d_{nx}}{2} \right)^2 \tau_L} \quad (3.1)$$

where I_{nx} is the effective laser intensity, E_{nx} is n % of the total energy on target, τ_L is the laser pulse length taken at FWHM and d_{nx} is the equivalent focal spot diameter containing n % of the total laser energy on target.

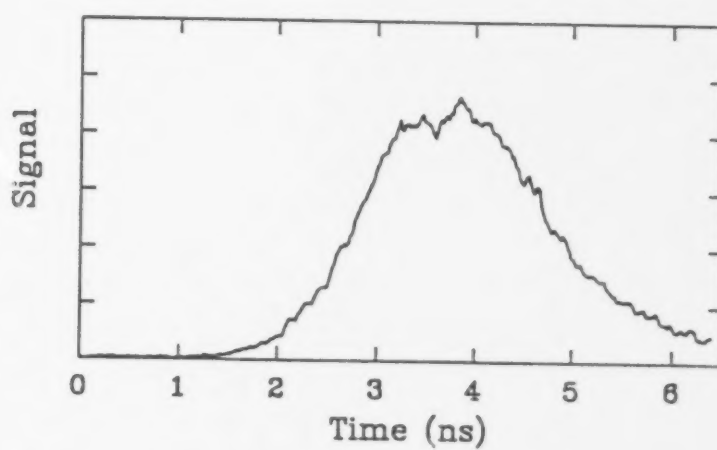
Table 3.1-3.4 summarizes the laser beam conditions on target for the three laser systems.



(a)

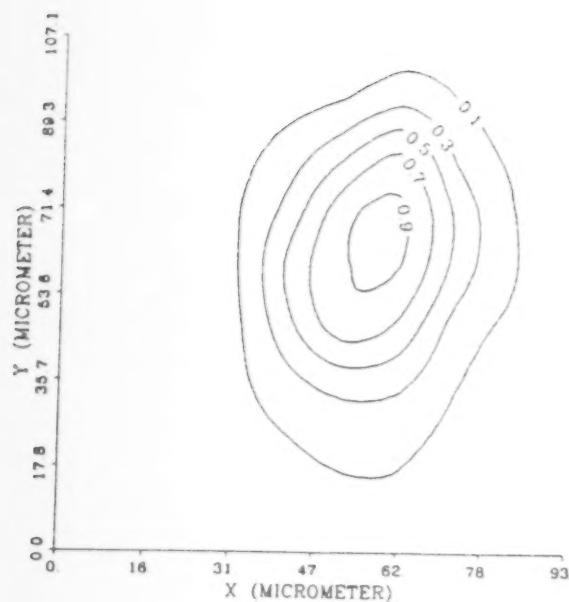


(b)

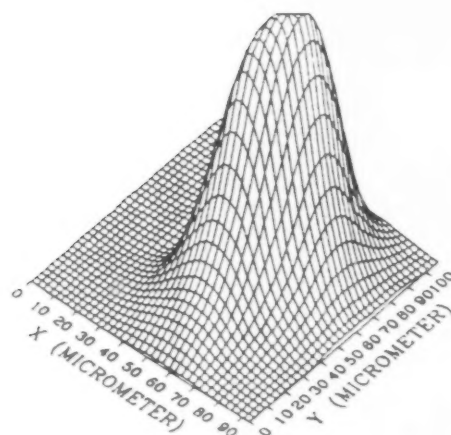


(c)

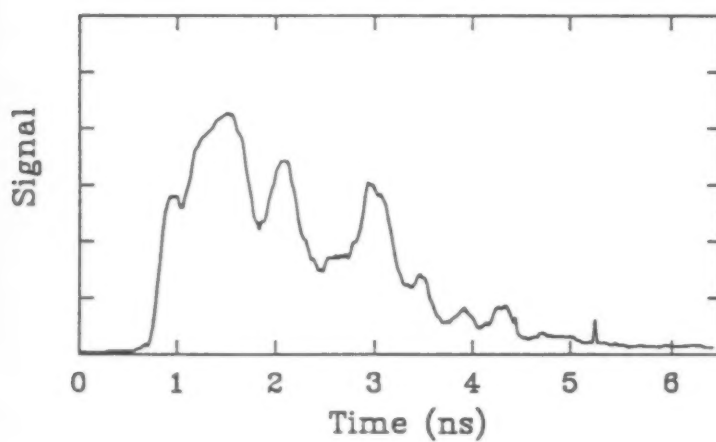
Figure 3.8 (a) Contour plot and (b) 3D mesh plot of the time integrated focal spot intensity distribution (c) temporal distribution of the Brillouin compressed KrF laser pulse.



(a)

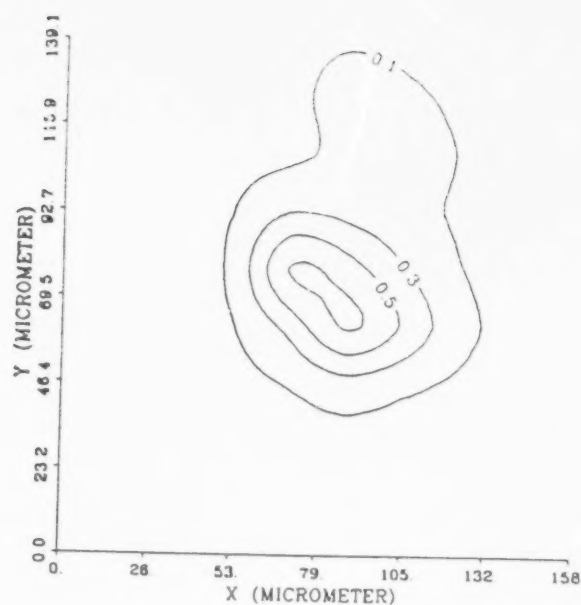


(b)

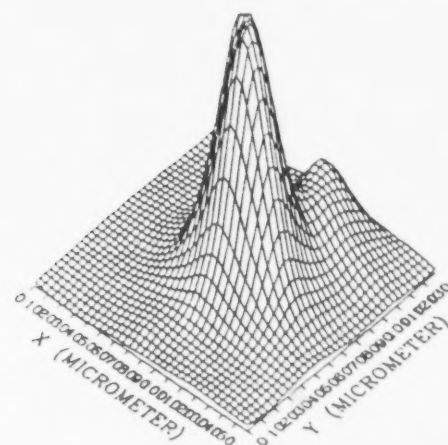


(c)

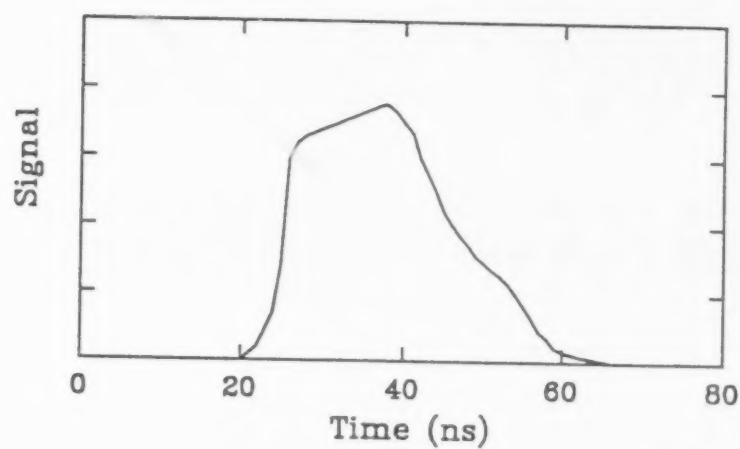
Figure 3.9 (a) Contour plot and (b) 3D mesh plot of the time integrated focal spot intensity distribution (c) temporal distribution of the Raman compressed KrF laser pulse.



(a)

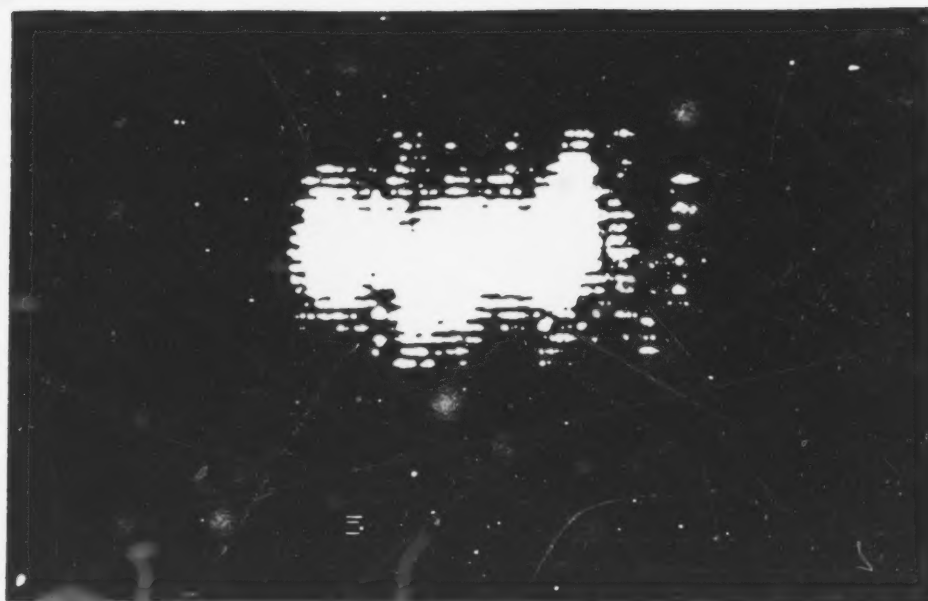


(b)

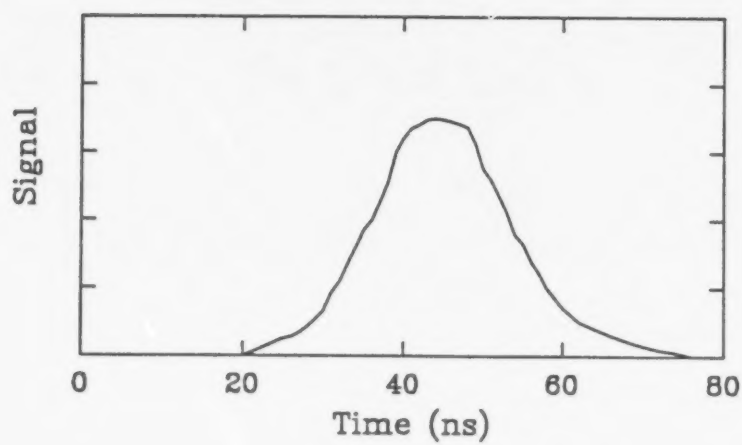


(c)

Figure 3.10 (a) Contour plot and (b) 3D mesh plot of the time integrated focal spot intensity distribution (c) temporal distribution of the uncompressed KrF laser pulse.



(a)



(b)

Figure 3.11 (a) Time integrated focal spot (b) temporal distribution of the ruby laser pulse.

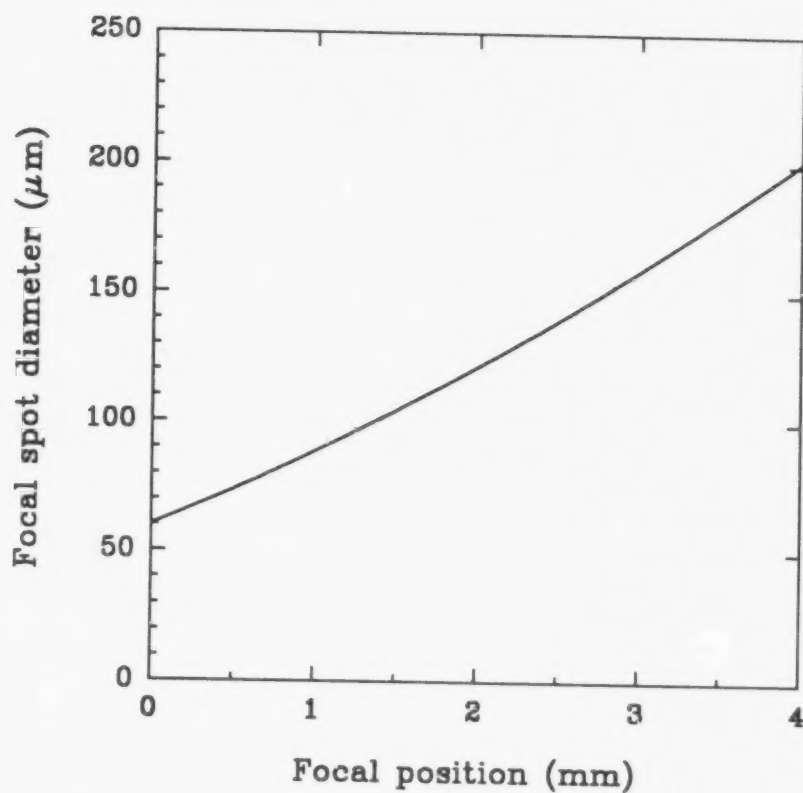


Figure 3.12 Equivalent focal spot diameter for 90 % energy contour as a function of the position of the focusing lens. For the Brillouin compressed KrF laser pulse.

Table 3.1 Raman Compressed KrF Laser Pulse Conditions

f	E (J)	τ (ns)	d90% (μm)	I90% (W/cm^2)	d50% (μm)	I50% (W/cm^2)
f0	1.0	2.0	74	1.0×10^{13}	34	2.8×10^{13}
f1	1.0	2.0	170	2.0×10^{12}	78	5.2×10^{12}
f2	1.0	2.0	340	5.0×10^{11}	156	1.3×10^{12}

Table 3.2 Brillouin Compressed KrF Laser Pulse Conditions

f	E (J)	τ (ns)	d90% (μm)	I90% (W/cm^2)	d50% (μm)	I50% (W/cm^2)
f0	1.0	2.0	68	1.2×10^{13}	28	4.1×10^{13}
f1	1.0	2.0	124	3.7×10^{12}	59	9.1×10^{12}
f2	1.0	2.0	203	1.4×10^{12}	102	3.1×10^{12}

Table 3.3 Front end KrF Laser Pulse Conditions

f	E (J)	τ (ns)	d90% (μm)	I90% (W/cm^2)	d50% (μm)	I50% (W/cm^2)
f0	0.5	20	119	2.0×10^{11}	28	7.0×10^{11}
f1	0.5	20	159	5.6×10^{10}	59	1.6×10^{11}
f2	0.5	20	203	3.7×10^{10}	102	1.0×10^{11}

Table 3.4 Ruby Laser Pulse Conditions

f	E (J)	τ (ns)	d90% (μm)	I90% (W/cm^2)	d50% (μm)	I50% (W/cm^2)
f0	0.5	25.0	100	2.3×10^{11}		

3.6 Targets Used in Experiments

Slabs or thick foils mounted on top of massive substrates of different materials are used as targets in the experiments. The materials are chosen to enhance different dominant physical processes. Carbon is used mainly for studying the recombination kinetics of an expanding plasma. Aluminum is used to study the energy transport of an intermediate-Z laser produced plasma. Gold is used to study the radiation transport of a high-Z laser produced plasma.

3.7 Integrating Sphere

An integrating sphere (Ulbricht sphere) consists of a hollow sphere with its inner surface coated by a diffuse reflector. The sphere (see figure 3.13) used in this thesis work was constructed by painting the inside of a 25.5 cm diameter copper sphere with an Eastman Kodak white reflective coating (cat.# 6080). Separate openings are provided for the incoming beam, the target holder and radiation detectors. Their overall area is less than 20% of the total surface area of the sphere. Radiation scattered from the target is multiply reflected by the diffuse reflector surface and monitored by four EG&G FND-100 ultra-fast silicon photodiode. They are covered by a quartz diffuser and a colored glass filter (UG5) to transmit only the laser light reflected by the sphere's inside wall. Absolute calibration of the sphere is carried out by illuminating directly the rear wall of the sphere with a low intensity laser pulse. The reflectivity of the target is then calculated by comparing the signal level of the photodiodes due to the light diffusely scattered from the target to that obtained previously from a calibration run.

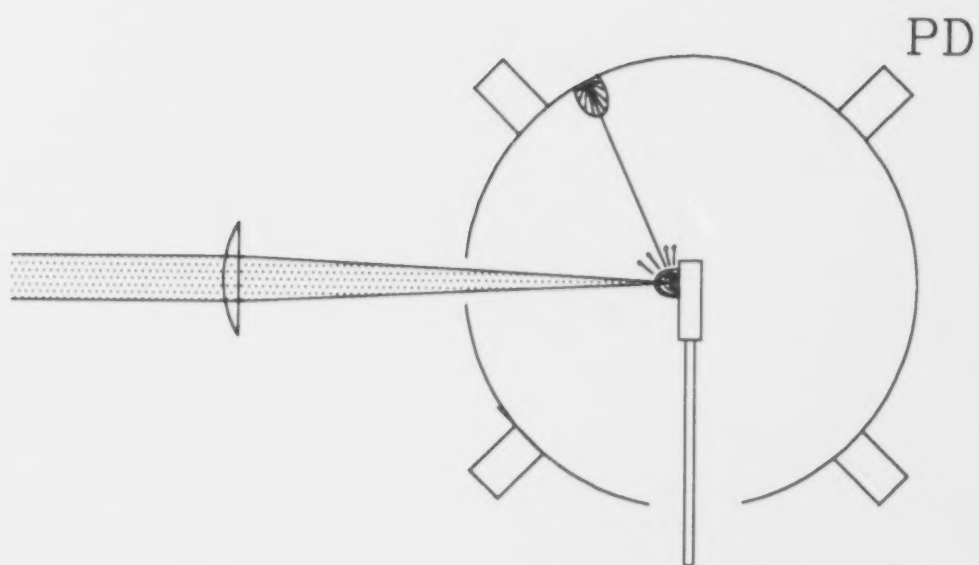


Figure 3.13 Schematic diagram of an integrating sphere.

3.8 X-Ray Foil Calorimeter

X-ray radiation from a laser produced plasma travels at the speed of light while the expanding plasma has a velocity typically less than 5×10^7 cm/s. At a distance of a few cm (say ~5cm) an energy detector with a time resolution of ~100 ns can discriminate between energy carried by x-rays and that carried by the expanding plasma.

Figure 3.14.a gives the circuit diagram of a fast response x-ray foil calorimeter developed in our laboratory²⁴. As detector element the device uses a free standing $2.04 \mu\text{m}$ thick aluminum foil mounted in the geometry shown in figure 3.14.b. Aluminum was chosen for its strength and availability in the desired thickness. A low atomic number absorber is preferred to one with a high atomic number to minimize secondary-electron generation^{25,26}. The thickness was chosen to ensure a fast time response, determined by the thermal diffusivity (for aluminum $0.97\text{cm}^2/\text{s}$). This gives a thermalization time of ~11 ns. The foil is free standing so that the cooling time due to lateral thermal conduction and radiation is very long (>ms). Thus the device effectively integrates the incident energy flux signal. When x-ray photons are absorbed by the foil, the temperature of the foil changes, which induces an increase in the resistance of the foil. A low-impedance wheatstone bridge with a pulsed voltage source is used to detect the changing resistance. The signal, V_s , is measured across the bridge circuit as shown in figure 3.14.a. A 50Ω coaxial cable is used to connect the bridge circuit to the measuring oscilloscope. This cable, which is terminated in a 50Ω load at its far end, appears as a 50Ω load resistor, R_L , in the circuit. The load resistance introduces only a minor perturbation on the bridge circuit since it is much higher in impedance than the foil element (1.0Ω). The complete circuit, together with the detector foil, was mounted in a small cast

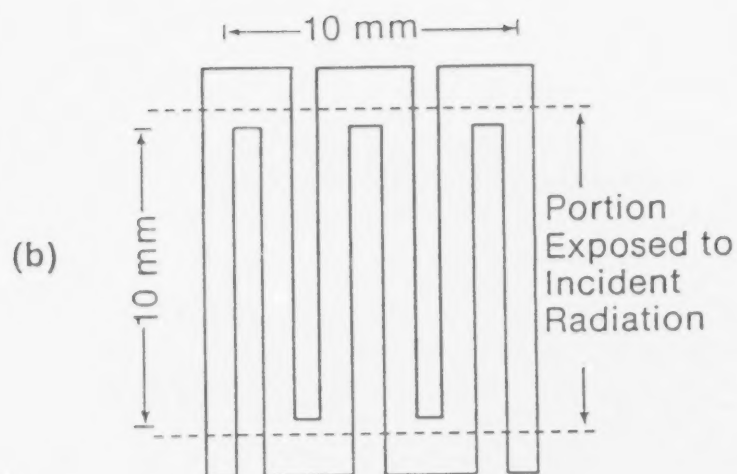
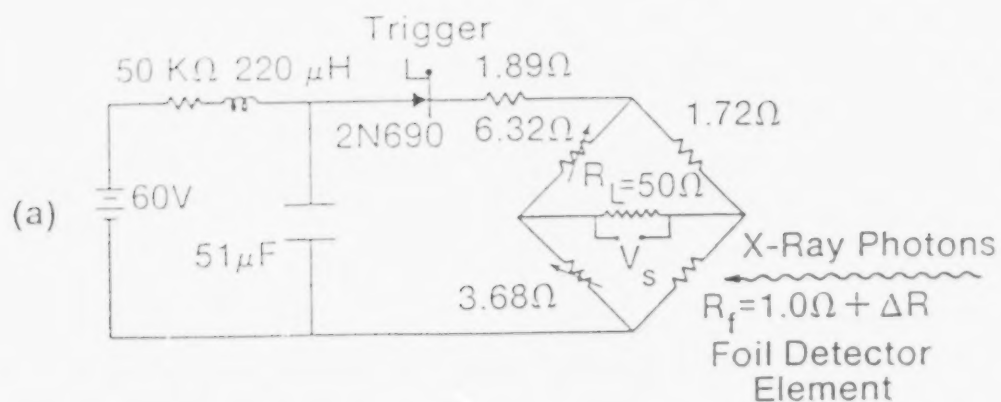


Figure 3.14 (a) The circuit diagram for the x-ray foil calorimeter. (b) A magnified diagram of the foil detector element.

aluminum box with an entrance hole for the x-ray radiation

The temperature rise experienced by the foil is 2.0°K for each 1 mJ/cm^2 absorbed energy. The corresponding fractional change in resistance is 8×10^{-3} times the foil resistance, R_f . For measurements reported in this thesis the absorbed energy flux was in the range of $0.2 - 5.0 \text{ mJ/cm}^2$. Given the small temperature rise, nonlinear effects due to the temperature dependence of the heat capacity and the change in temperature coefficient for resistivity are negligible. Similarly, nonlinear second order terms in the circuit response are negligible for the variation in R_f .

Calibration of detector sensitivity was carried out by using a 20 ns KrF laser pulse with an energy flux of $\sim 30 \text{ mJ/cm}^2$. The fraction of KrF laser radiation absorbed by the $2 \text{ }\mu\text{m}$ aluminum foil was determined by using a separate but similar piece of sample foil mounted inside an integrating sphere, to measure the fraction of non-absorbed energy. In this manner, the response function was determined to be $41.5 \pm 5 \text{ V cm}^2/\text{J}$. This value is in excellent agreement with the nominally calculated value of $43.7 \text{ V cm}^2/\text{J}$ for the values of the circuit elements.

When the device is used for measuring the x-ray yield from laser-produced plasma the ions from the expanding plasma arrive at the detector and short out the bridge circuit, leading to excessively large signals on the measuring oscilloscope. To reduce these plasma effects in the majority of cases, a magnetic field was used across the entrance of the calorimeter box. Using a permanent magnet, a field of 850 gauss was applied perpendicular to the trajectory of the ions. Figure 3.15 gives a typical output from the device.

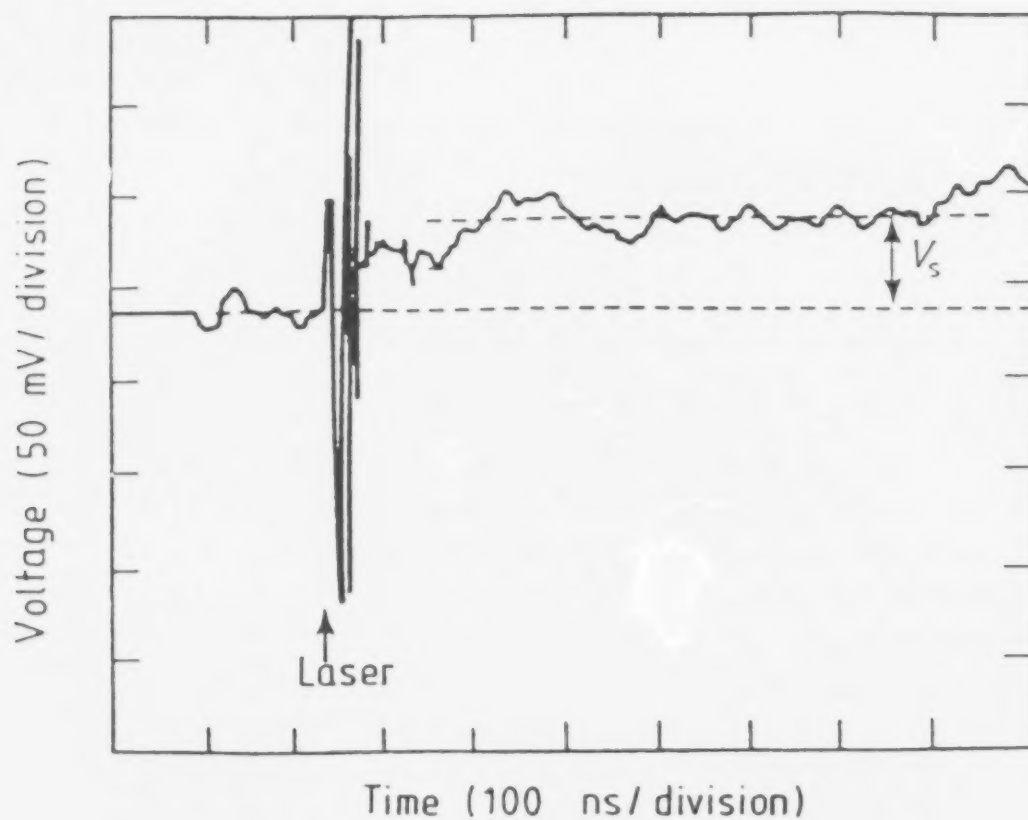


Figure 3.15 The output signal of a x-ray foil calorimeter for a Au target irradiated at a laser intensity of $5 \times 10^{12} \text{ W cm}^2$, with a magnetic field to reduce the detection of ions.

3.9 Plasma Calorimeter

Plasma calorimeters are used to determine the energy in the expanding plasma and emitted by x-rays. Plasma and x-ray energy absorbed in the bulk of a massive material results in a change of temperature of the material. Two types of plasma calorimeters are used in this thesis work. In the first type an aluminum plate is used as the energy absorber and a commercial Peltier effect thermoelectric device is used to measure the temperature change. The sensitivity of the device is of order of a few mV/J and signal amplification is therefore required. The device can be used in a differential configuration to discriminate against the scattered light by using a filter placed in front of one of the detectors. Absolute calibration is done in situ by electrically depositing a known amount of energy into a calibration resistor imbedded inside the aluminum plate absorber. The second type of plasma calorimeter is similar to the first except that the aluminum plate absorber is replaced by a thin anodized aluminum foil. Calibration is done optically by irradiated the device with a known amount of laser energy.

3.10 Foil/Scintillator/Photomultiplier Detectors

The temperature monitor consists of four channels each equipped with a metal filter (usually aluminum) in front of a 3 mm thick NE 102²⁷ plastic scintillator coupled to a Hamamatsu R1450 photomultiplier via a fiber optic bundle. The light input to the photomultiplier is attenuated by neutral density filters, when required, to ensure linearity. The output of the photomultipliers is recorded by 4 channels of a LeCroy 2249A ADC charge integrator. Calibration of the monitor is performed by using identical aluminum filters. The relative response of a channel is determined by including the spectral emission function of

the plastic scintillator coupled with the spectral response functions of the photomultiplier and the neutral density filters (see figure 3.16). During normal operation, aluminum filters of thickness of approximately 5 μm , 9 μm , 25 μm and 50 μm are used. For a 5 μm aluminum filter the cut off energy is approximately 1 keV. The electron temperature (T_e) from the plasma is determined by the foil transmission technique²⁸. For carbon and aluminum plasmas, estimation of T_e is obtained by fitting the experimentally observed relative x-ray intensities for different foils with those calculated using a collisional-radiative (CR) model²⁹ (see figure 3.17). The ion number density (n_i) is assumed to be between 10^{20} cm^{-3} 10^{21} cm^{-3} . For our experimental condition, and aluminum filters of thickness greater than few μm ($\sim 3 \mu\text{m}$), the x-ray line radiation is effectively attenuated leaving primarily the continuum radiation. The relative x-ray emission is insensitive to the electron number density of the plasma. In the case of gold plasma, estimation of T_e is obtained by fitting the experimentally observed relative x-ray intensities for different foils with those calculated using a coronal model³⁰.

3.11 Charge Collectors

Charge collectors or Faraday cups^{31,32,33} measure the total current arriving at the detector as a function of time

$$I(t) = T n_i \langle Z \rangle e v \quad (3.2)$$

where n_i is the number density of ions, $\langle Z \rangle e$ is the average charge of ions of velocity v , and T is a geometric factor depending only on the collector. A schematic diagram of a charge collector is shown in figure 3.18.a. The main element consists of a biased entrance grid to extract ions

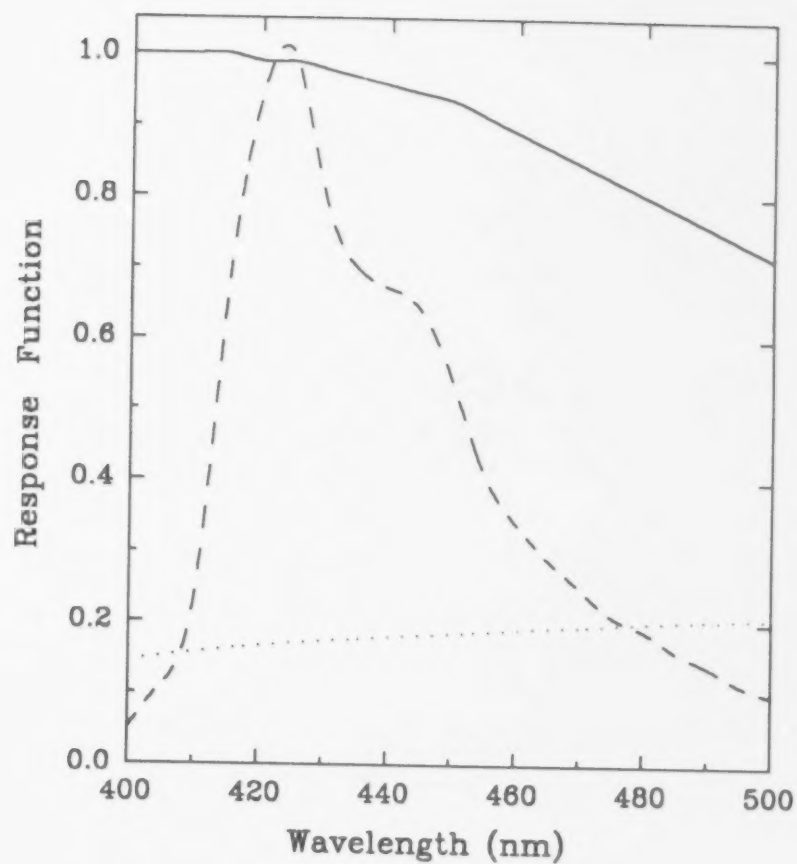


Figure 3.16 The spectral emission function of the scintillator (dashed line) and the spectral response functions of the photomultiplier (solid line) and neutral density filter ND.6 (dotted line).

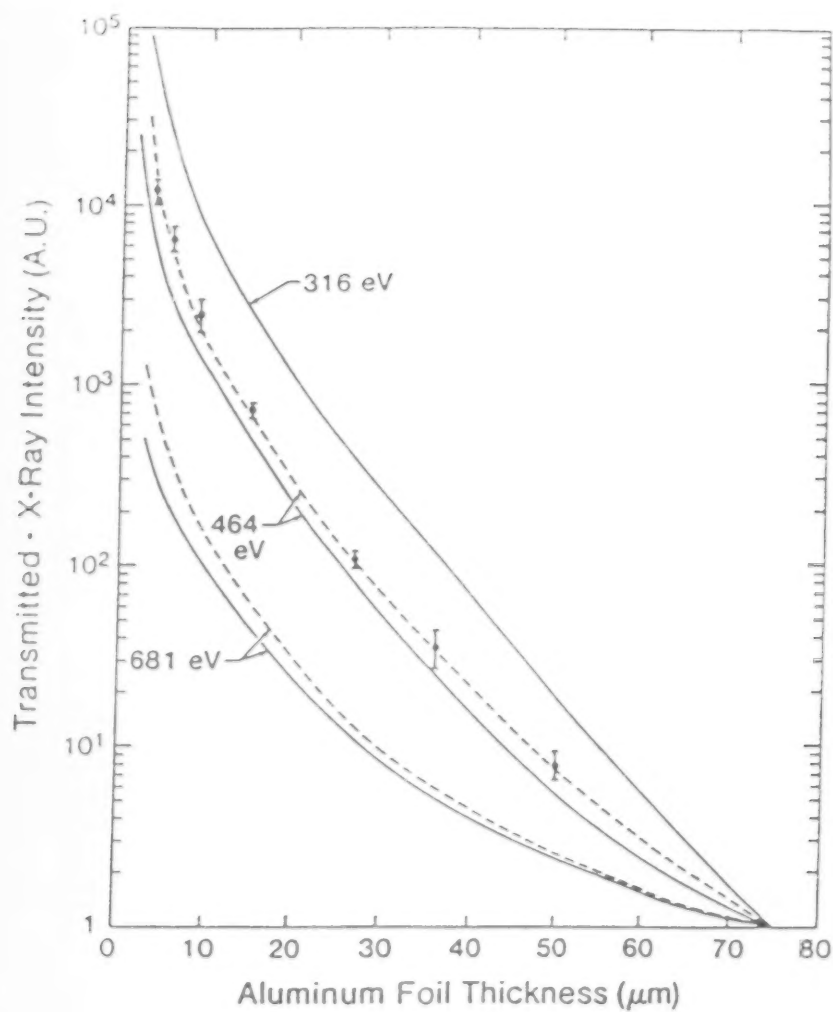


Figure 3.17 Variation of the x-ray intensity transmitted through aluminum foils computed on the basis of CR model: $n_i = 10^{21} \text{ cm}^{-3}$ (solid line), $n_i = 10^{20} \text{ cm}^{-3}$ (dotted line). Full circles represent the experimental data.

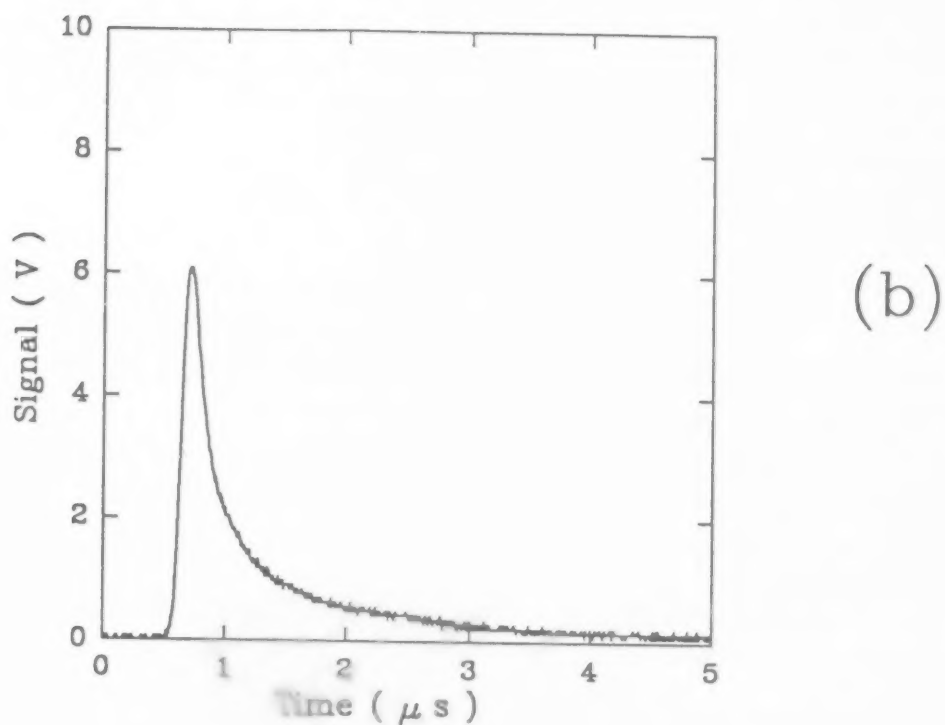
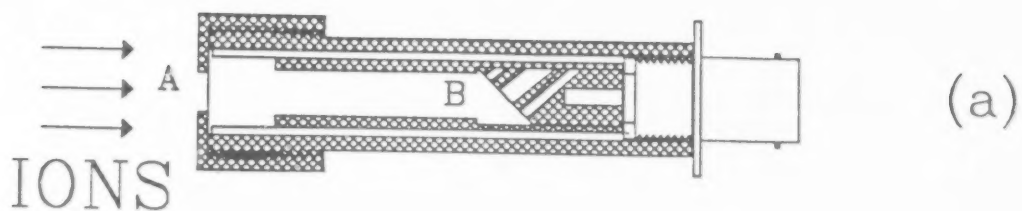


Figure 3.18 (a) A schematic diagram of a Faraday cup: A, entrance grid; B, collector cup (b) Typical output from a Faraday cup at 25 cm from the target surface. The carbon ions are produced by a $9 \times 10^{12} \text{ W/cm}^2$ KrF pulse.

from the plasma and a deep collector cup to re-collect the secondary electrons produced by the impact of the ions on the cup surface. A blocking filter is placed at the output to separate the bias voltage from the output signal. The extraction of ions is done by using a potential on the grid surface that is sufficiently negative relative to the plasma. Since the potential decreases as $-\exp(-d/\lambda_D)$ at a distance d from a charged surface, the grid spacing must be less than the Debye length λ_D to ensure the grid forms a uniform potential surface. This device has a linear response up to $-10V$, beyond which space charge effects will lead to output saturation. A typical output signal of the device is shown in figure 3.18.b.

3.12 Charge Analyzers

A charge analyzer is a device that can separate a multi-charge ion beam into separate single-charge ion beams. An ion beam is extracted from the expanding plasma by stripping its electrons. An electric field is used to separate the ion beam into individual ion beams. The governing equations for the trajectory of the ion beams are

$$\frac{dx}{dt} = \frac{x}{t} \quad (3.3)$$

$$\frac{d^2y}{dt^2} = \frac{Z e}{m} E(x,t) \quad (3.4)$$

where the ions initially propagate along the $\hat{x}$ direction and $E(x,t)\hat{y}$ is the applied electric field. Two types of charge analyzers are used in this thesis work as described below.

3.12.1 Electrostatic Charge Analyzer

A parallel plate electrostatic charge analyzer^{34,35} was first used in this thesis work to detect ions of different charge state. Figure 3.19 gives a schematic diagram of the device used in this work. It contains two capacitor plates separated by a distance d with one at ground and the other at positive high voltage. The plate which contains the entrance slit is grounded in order to screen the applied field from the electrical noise of the plasmas. The device also contains a guard plate to prevent too much plasma entering the analyzer that could introduce high levels of unwanted electrical noise or even electrical break down. Guard rings are placed in between the two plates to minimize fringe field nonuniformities and perturbation of the electrostatic field due to the incoming ions. The charge collectors are in a deep cup configuration and biased at a sufficiently positive voltage (~ 50 V) to prevent escape of secondary electrons produced by ion impact on the collector surface.

The device separates ions of different charge state according to their kinetic energy to charge state ratio (e/Z) and time of flight selection. Ions of the same e/Z will have the same trajectory when deflected by a constant electric field.

$$x = \frac{2}{e} \frac{e}{Z} \frac{d}{V} \sin(2\theta) \quad (3.5)$$

However, the time of arrival of these ions at the detectors will depend on their kinetic energies.

$$t = \sqrt{\frac{2 \frac{e}{Z} \sin(2\theta)}{m_i \left(d + \frac{x}{\cos(\theta)} \right)}} \quad (3.6)$$

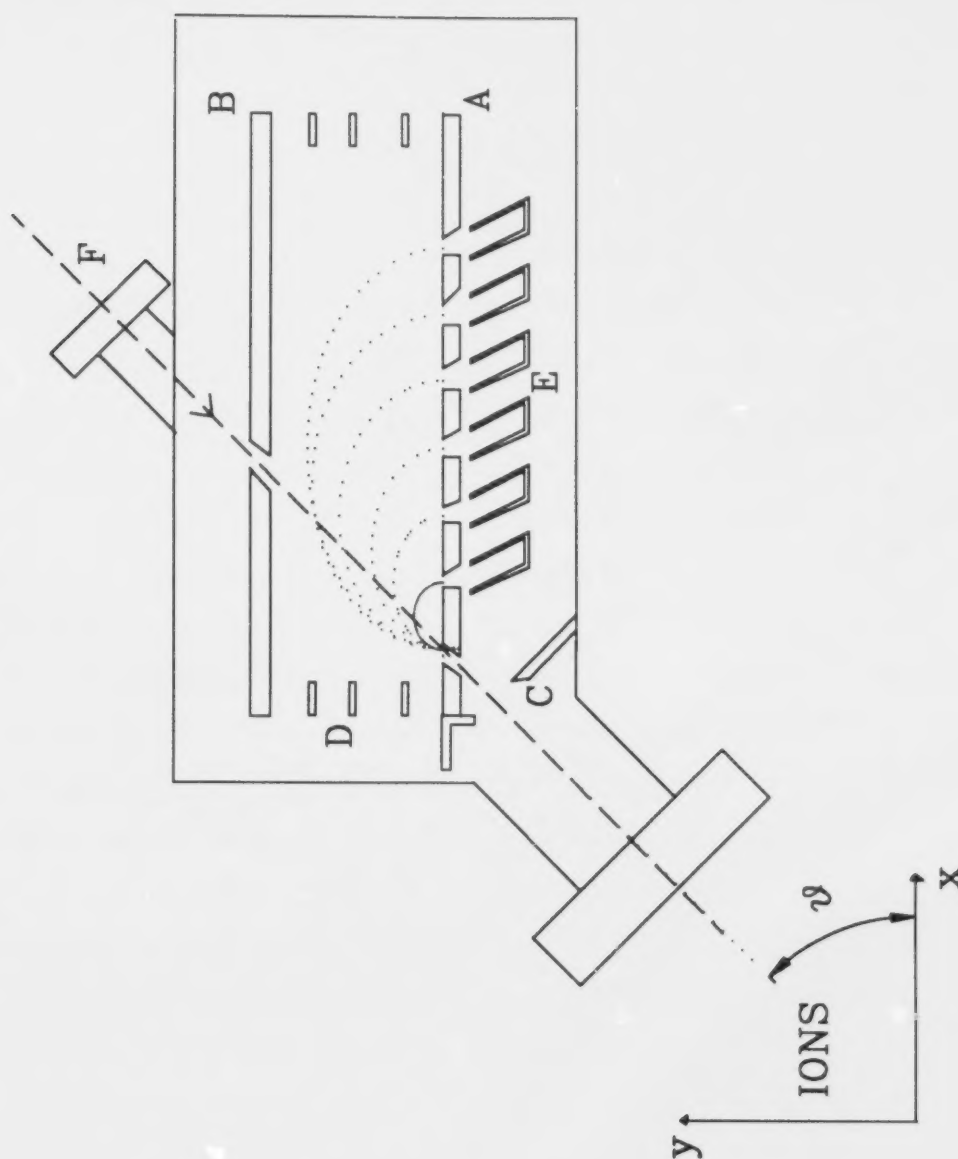


Figure 3.19 The schematic diagram of an electrostatic charge analyzer: A, ground plate; B high voltage plate; C, guard plate; D, guard rings; E, collector cups; F, alignment Laser.

Thus ion species having the same ϵ/Z but different Z are recorded by the same channel as temporally spaced pulses. A device which has the entrance and exit slits on the same plane and $\theta = 45^\circ$ possesses a first order focusing property. This can be demonstrated by using a Taylor expansion for δx in terms of $\delta\theta$

$$\delta x = \left(\frac{dx}{d\theta} \right) \delta\theta + \frac{1}{2} \left(\frac{d^2x}{d\theta^2} \right) \delta\theta^2 + O(\delta\theta^3) \quad (3.7)$$

For $\theta = 45^\circ$, $dx/d\theta$ vanishes and $\delta x = O(\delta\theta^2)$. The energy resolution of a given analyzer is usually governed by the finite lifetime and dimension of the ion source and the deviation of beam angle $\delta\theta$ due to finite dimension of the entrance slit. In laser produced plasma experiments the effect due to $\delta\theta$ dominates, implying

$$\frac{\delta\epsilon}{\epsilon} \approx 2 \delta\theta \quad (3.8)$$

Because of the self focusing properties of the device, the longitudinal space charge expansion does not result in a loss of charge. However, it changes the ion beam paths resulting in an increase of charge collection time such that the pulse resolution is degraded and output current is decreased³⁵. This implies the existence of a cut off velocity beyond which the desired time resolution is no longer satisfied.

The entrance and exit slits with dimensions L_{ent} by W_{ent} and L_j by W_j respectively are cut at 45° (i.e. $\theta = 45^\circ$) to the y-axis. The parameters L_{ent} , W_{ent} , x_j , L_j , W_j , and d must be chosen carefully. Space charge effects within a charge analyzer can cause saturation of output signal for an entrance slit wider than some maximum value³⁵. For a given slit length, usually chosen to be around a centimeter, this value can be determined experimentally³⁵. The minimum value of the slit width is determined by the

output signal to noise ratio (SNR) which is usually required to be ten. The output slit dimension should be chosen large enough to avoid loss in signal. It is desirable to choose $x_{j+1} = \alpha x_j$ where α is a constant. The relative calibration between channels can, in principle, be obtained in 2 shots by changing the applied voltage a factor of α from each other, provided the two laser pulses are identical. The change of voltage causes the channel which samples a given ϵ/Z to shift over by one. Finally, the deflection distance d should be chosen by the condition that ions to be detected must not strike the upper plate. Figure 3.20 show a typical output pulse from the charge analyzer.

3.12.2 Electrodynamic Charge Analyzer

In order to give simpler charge state temporal histories on a single shot basis an electrodynamic charge analyzer^{36,37,38} was developed for this thesis work. As described in section 3.12.1, ions deflected by a constant electric field will have trajectories that depend on their kinetic energy to charge state ratio. Since the kinetic energy of an ion is proportional to $1/t^2$, ions deflected by an electric field that varying according to $1/t^2$ will have the same trajectories.

Fig 3.21 show a schematic diagram of the device. The entrance consists of two pinholes 5 cm apart to limit the divergence of the input ion beam. A pair of parallel plate electrodes are used to separate the ion beam into beams of different charge states. The charge state separated ion beams are collected by the charge collector cups in a deep cup configuration to suppress secondary electron emission. The size of the second pinhole is chosen to be 200 μm to limit the density of the ion beam so that the space charge spread of the ion beams at the collector cups is less than the entrance aperture of the cups. A y-configuration

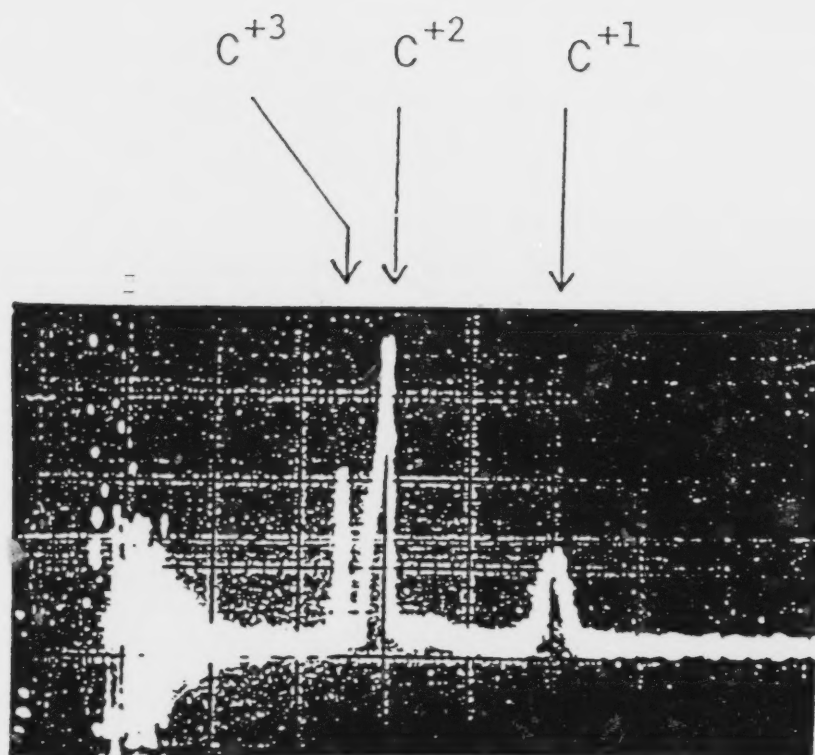


Figure 3.20 Typical output from an electrostatic charge analyzer at a distance 40 cm from the target surface. The carbon ions are produced by a $1 \times 10^{11} \text{ W/cm}^2$ ruby laser pulse. (50 mV, 1 μs per division)

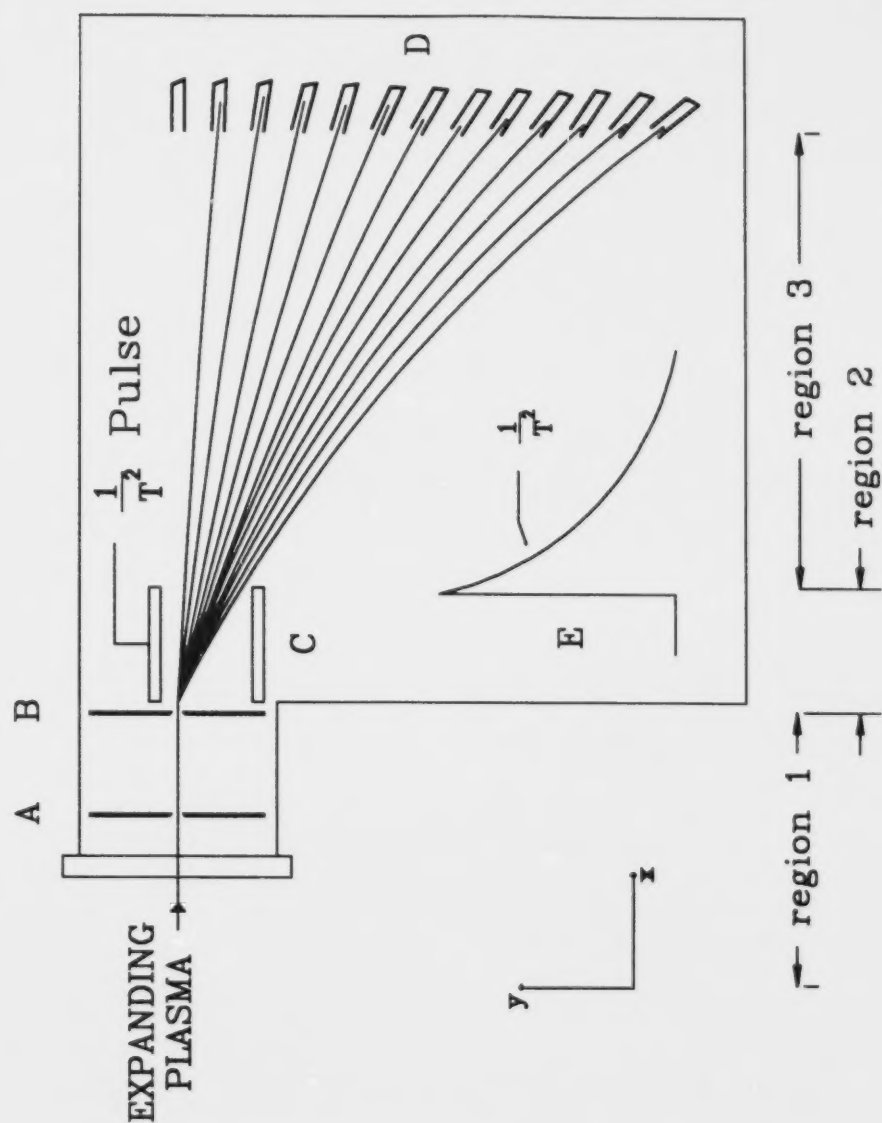


Figure 3.21 Schematic diagram of an electrodynamic charge analyzer. A, entrance aperture; B, alignment aperture; C, deflection plates; D, charge collector cups; E, electric field across the deflection plates.

(deflection in the y-direction) is chosen in this design to reduce the voltage requirement of the analyzer circuit using a $1/t^2$ electric field starting at a time τ after firing the laser, described by

$$E(x,t) = \begin{cases} 0 & \text{region I} \\ \frac{V(t)}{d} & \text{region II} \\ 0 & \text{region III} \end{cases} \quad (3.9)$$

where

$$V(t) = \frac{V_0 \tau^2}{t^2} \quad (3.10)$$

From equations (3.3) and (3.4) the deflection distance of an ion is given by

$$Y_3 = \frac{Z e V_0 \tau^2}{m d} \left(x_3 \left(\frac{1}{x_1} - \frac{1}{x_2} \right) - \ln \left(\frac{x_2}{x_1} \right) \right) \quad (3.11)$$

where x_1 is the distance from the target surface to the entrance of the charge analyzer; x_2 is x_1 plus the length of the region where high voltage deflection field is applied and x_3 is x_1 and x_2 plus the length of the field free region. The $1/t^2$ voltage pulse can be generated by using a resistor-capacitor passive network which transforms a square pulse into the desired pulse shape. The function can be approximated by an exponential series

$$f(t) = \sum_j A_j \exp(a_j t) \quad (3.12)$$

A Laplace Transform of the exponential series is an algebraic function which can be realized as a RC driving point function³⁹. A RC driving point function can be represented by Foster's second form that can be generated

in analog form by a passive network consisting of a capacitor, a resistor, and a series of RC combinations all in parallel. Figure 3.22.a illustrates a circuit which is capable of transforming a step voltage into pulse which in theory approximates the $1/t^2$ pulse to within 1%, with delay time τ at 5 μ s. Figure 3.22.b compares the output of the circuit and the ideal $1/t^2$ pulse. The value of the required step voltage can be calculated from equation (3.11). Circuits for different delay times can be obtained by scaling the RC time constant of each ladder linearly. For example, with a 1 μ s delay time the resistances of the circuit can be kept unchanged and the capacitances of the circuit decreased by a factor of five. The dimension of the entrance pinholes and exit slits must be chosen in such a way that there is no loss of signal due to the spread caused by the space charge field. The space charge field can cause an axially symmetric ion beam to expand according to the following relation³⁸

$$\int_0^{y_0} \exp(y^2) dy = \sqrt{\frac{Z e J d^2}{m_1 4 \epsilon_0 v^3}} \quad (3.13)$$

where

$$y_0 = \sqrt{\ln \left(\frac{R}{r_0} \right)} \quad (3.14)$$

where r_0 is the radius of the beam at $x = 0$ and R is the expanded radius at $x = d$. J is the average current density of the ion beam at $x = 0$, and v is the axial velocity of the ions. The required exit slit can be chosen based on equation (3.13) and (3.14). To ensure the charge collector cups capture all the signal the spread of the ion beams was experimentally measured by replacing the collector cup with an array of small plates 1 mm wide and 1/2 mm apart.

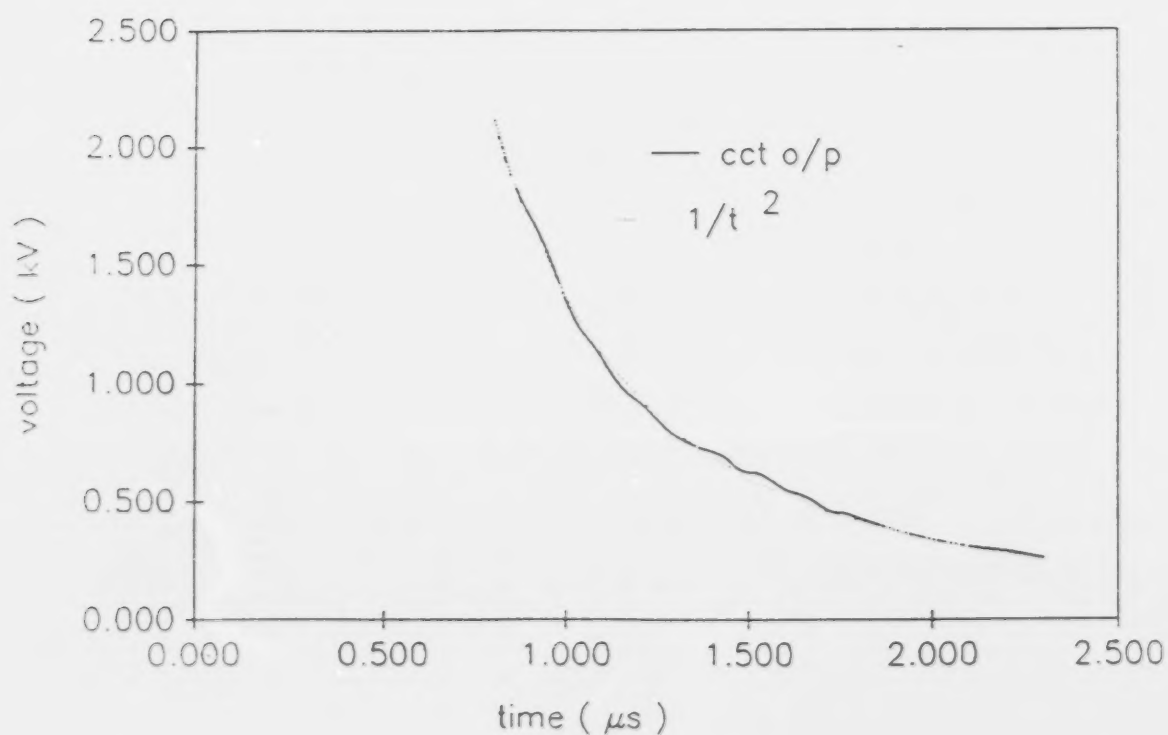
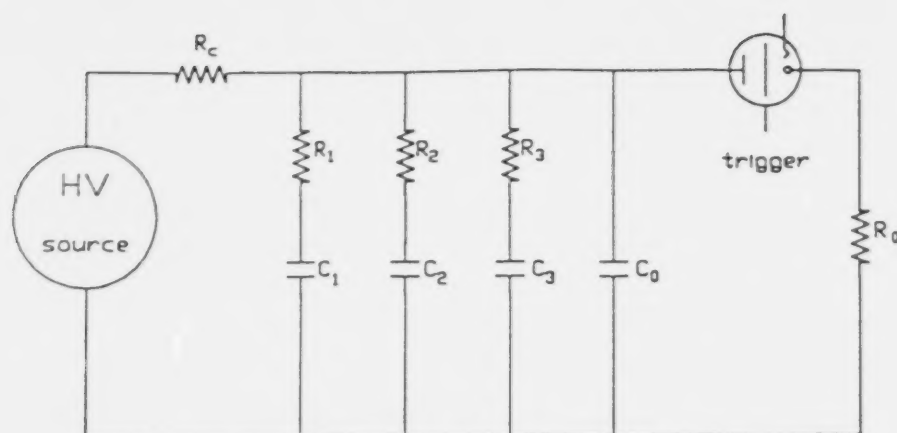


Figure 3.22 (a) Circuit diagram for $1/t^2$ circuit (b) Output from an $1/t^2$ circuit, real circuit output (solid) and ideal $1/t^2$ curve.

The final charge analyzer was set up such that the charge states of carbon C^{1+} to C^{6+} should appear on channels 2, 4, 6, 8, 10, 12 respectively and the intermediate channels should show no signal. Figure 3.23 illustrates that a complete charge state spectrum can be obtained for each shot thus relaxing the requirement on reproducibility of laser shots. In this case no C^{5+} or C^{6+} ions were present. The data shown represents the raw data without using any averaging process. As can be seen, the electrical noise level is significant and generally an 11 point gaussian averaging is used to smooth the data somewhat.

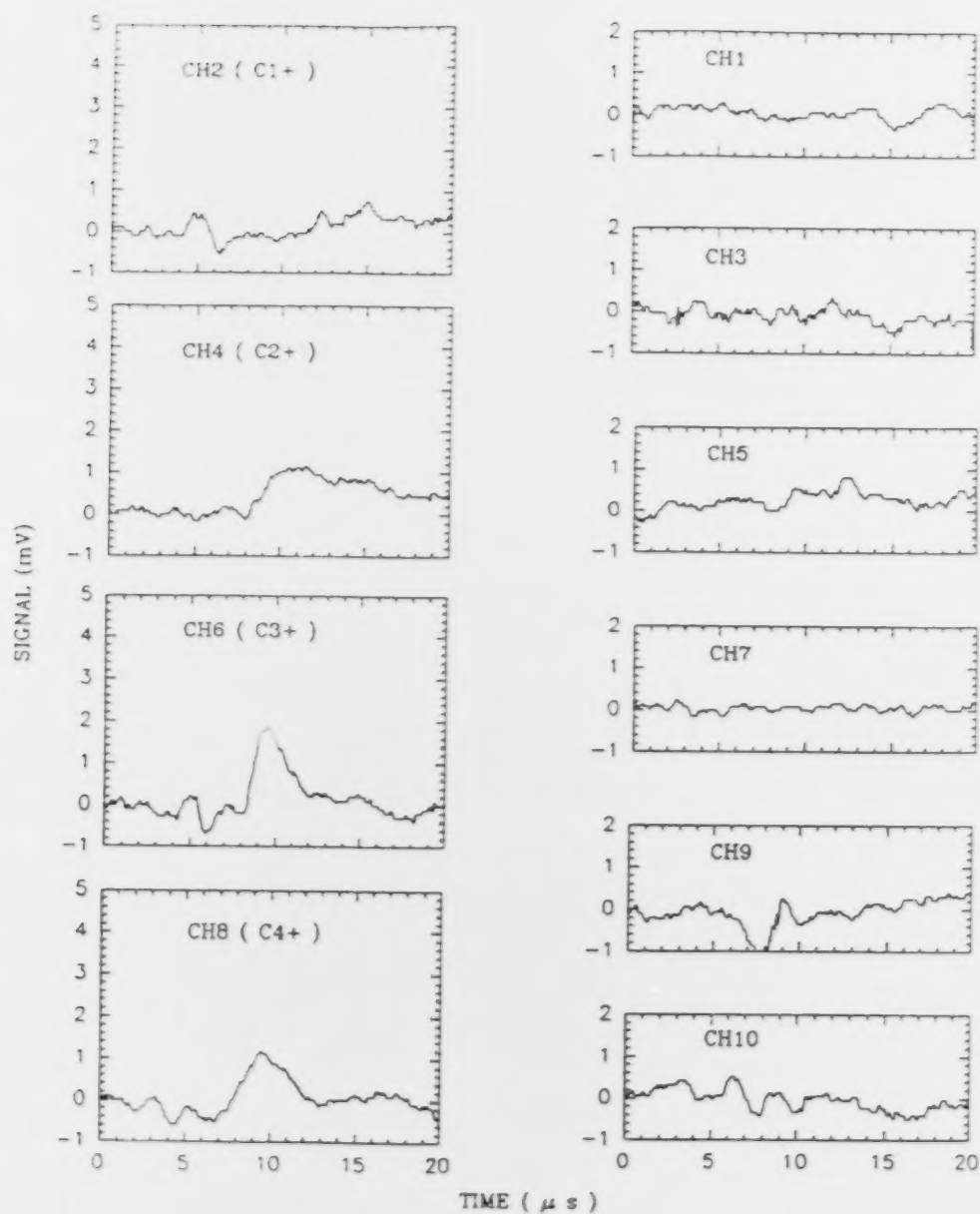


Figure 3.23 Typical output from an electrodynamic charge analyzer at a distance 80 cm from the target surface. The carbon ions are produced by a $3 \times 10^{11} \text{ W/cm}^2$ ruby laser pulse.

CHAPTER FOUR

ASYMPTOTIC BEHAVIOR OF THE EXPANDING PLASMAS

The study of ion expansion characteristics in laser-produced plasma plays an important role in understanding the physical processes involved in ablation⁴⁰⁻⁴³ and ion recombination^{44,45} and in determining ablation parameters⁴⁶⁻⁴⁹.

The ion expansion characteristics of carbon, aluminum and gold plasma created by focusing a laser beam on planar targets in the intensity range of $10^{10} - 10^{13} \text{ W/cm}^2$ will be reported in this chapter. The laser parameters (see section 3.5) consist of laser wavelength ranging from $0.25 \mu\text{m}$ to $0.69 \mu\text{m}$, pulse length ranging from 1 ns to 25 ns and the equivalent 90% energy containing focal diameter ranging from $68 \mu\text{m}$ to $340 \mu\text{m}$.

Arrays of plasma calorimeters and Faraday cups positioned at 10 - 25 cm from the target surface were used to map the angular distribution of energy and ion velocity in the expanding plasma (see figure 4.1). The energy measurements will be reported in chapter five.

The velocity distribution of ions at the ion detectors is essentially governed by the hydrodynamic expansion of the hot plasma into vacuum where the initial thermal energy is converted to kinetic energy. However, recombination occurring during expansion alters the charge state of ions. Knowledge of the variation of the charge state of ions is required to determine the velocity distribution of ions from the ion current traces.

An electrostatic charge analyzer and an electrodynamic charge analyzer positioned at 60 - 80 cm from the target surface were used to obtain charge state distribution of the expanding plasmas.

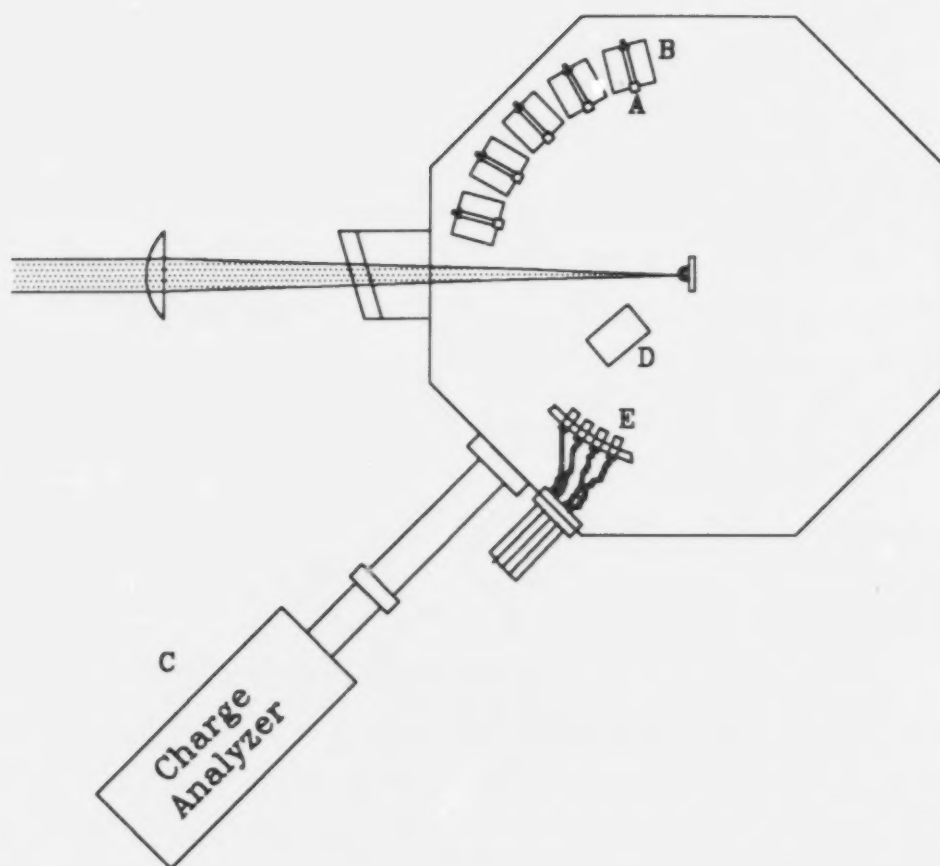


Figure 4.1 Experimental set up: arrays of charge collectors (A); arrays of plasma calorimeters (B); charge analyzers (C); x-ray calorimeter (D) and x-ray absorbing foil scintillator detectors (E).

4.1 Ion Expansion Characteristics of Carbon Plasma

Faraday cup and charge analyzer measurements of ions from an expanding plasma produced by focusing a 2 ns, 0.248 μm laser beam on planar carbon targets in the intensity range of $10^{10} - 10^{13} \text{ W/cm}^2$ are reported in section 4.1.a and 4.1.b, those for a 20 ns, 0.248 μm laser beam with a nominal laser intensity value of $2 \times 10^{11} \text{ W/cm}^2$ are reported in section 4.1.c and 4.1.d and those for a 25 ns, 0.694 μm laser beam with a nominal laser intensity value of $1 \times 10^{11} \text{ W/cm}^2$ are reported in section 4.1.e and 4.1.f.

4.1.a Ion Emission Measurements for Short Pulse KrF Laser Produced Carbon Plasma

The ion current traces observed in a direction $\theta = 14^\circ$, 28° , 42° and 56° to the target normal for different laser intensities are shown in figures 4.2 to 4.5. For a fixed direction of observation, it is found that the velocity of the ion peak increases with increasing laser intensity. As the focusing condition is changed from best focus to 4 mm out of focus, the FWHM of the ion current traces become broader. As the direction of observation is changed, the magnitude and velocity of the ion peak decreases with increasing value of θ .

The laser intensity scaling of the velocity of the peak of the ion current traces observed in $\theta = 14^\circ$, 28° , 42° and 56° are shown in figure 4.6. The best fit to these data using a simple power law curve can be summarized in the following equation

$$v_{pk} = A \times 10^7 I_{13}^B \quad [\text{cm/s}] \quad (4.1)$$

where v_{pk} is the velocity of the peak of the ion currents and I_{13} the laser intensity in units of 10^{13} W/cm^2 ; and A

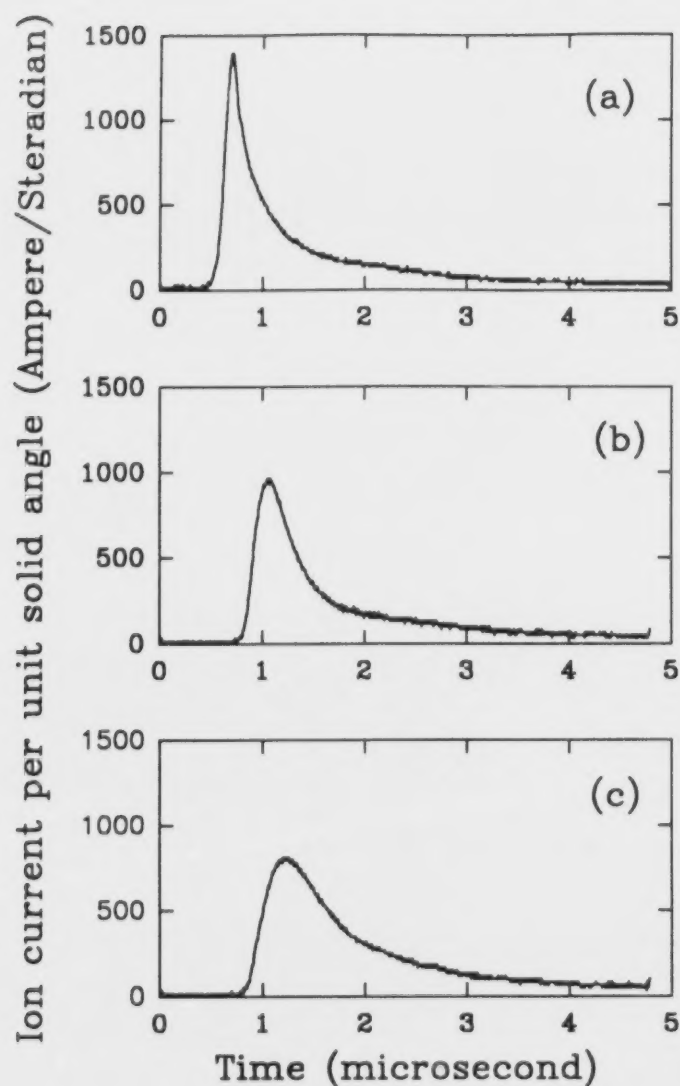


Figure 4.2 Ion current signals observed by a Faraday cup at 25 cm from the target surface and 14° to the target normal. The C plasmas were produced by laser pulses of $0.248 \mu\text{m}$ and 2 ns at irradiance of (a) $I_L = 9.1 \times 10^{12} \text{ W/cm}^2$ (best-focus); (b) $I_L = 2.7 \times 10^{12} \text{ W/cm}^2$ (2 mm-off-focus); and (c) $I_L = 8.7 \times 10^{11} \text{ W/cm}^2$ (4 mm-off-focus).

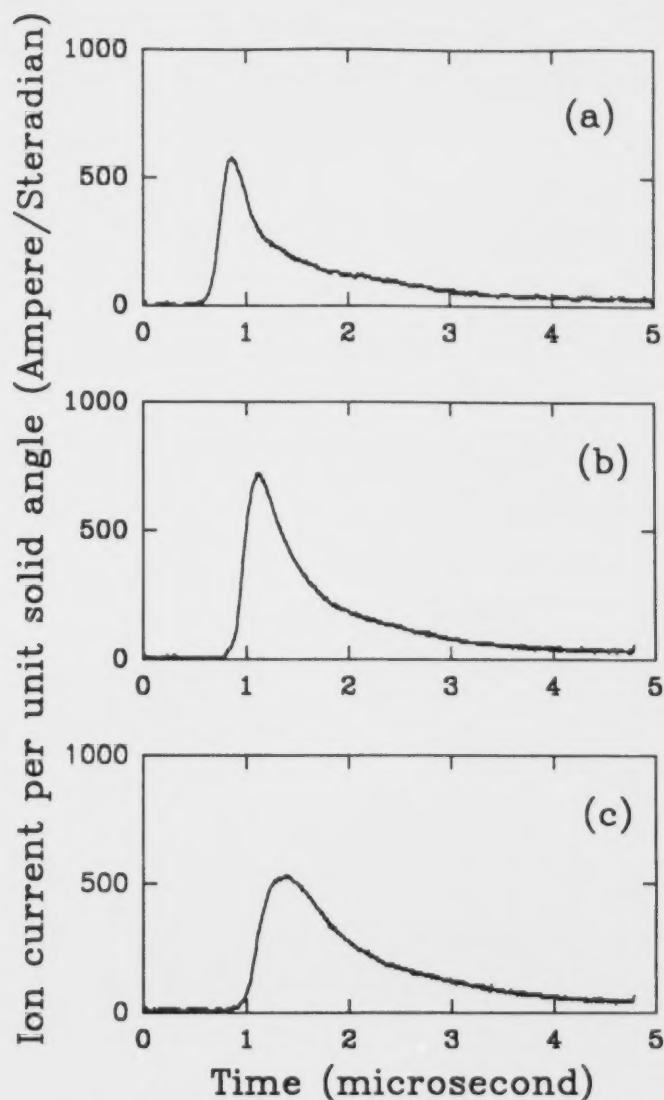


Figure 4.3 Ion current signals observed by a Faraday cup at 25 cm from the target surface and 28° to the target normal. The C plasmas were produced by laser pulses of $0.248 \mu\text{m}$ and 2 ns at irradiance of (a) $I_L = 9.1 \times 10^{12} \text{ W/cm}^2$ (best-focus); (b) $I_L = 2.7 \times 10^{12} \text{ W/cm}^2$ (2 mm-off-focus); and (c) $I_L = 8.7 \times 10^{11} \text{ W/cm}^2$ (4 mm-off-focus).

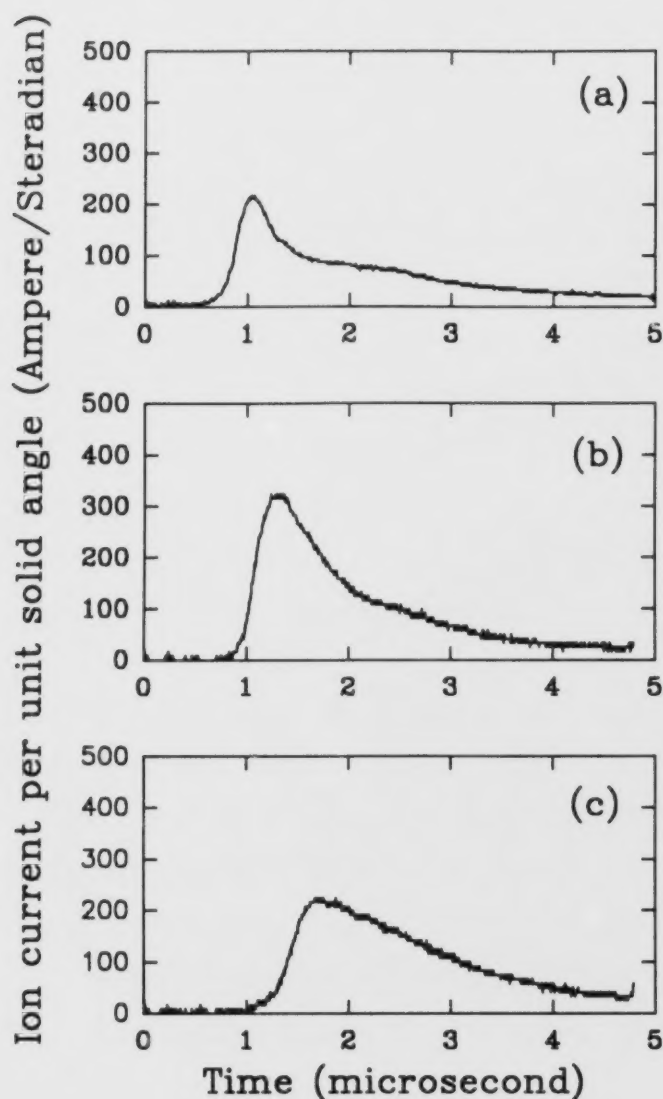


Figure 4.4 Ion current signals observed by a Faraday cup at 25 cm from the target surface and 42° to the target normal. The C plasmas were produced by laser pulses of $0.248 \mu\text{m}$ and 2 ns at irradiance of (a) $I_L = 9.1 \times 10^{12} \text{ W/cm}^2$ (best-focus); (b) $I_L = 2.7 \times 10^{12} \text{ W/cm}^2$ (2 mm-off-focus); and (c) $I_L = 8.7 \times 10^{11} \text{ W/cm}^2$ (4 mm-off-focus).

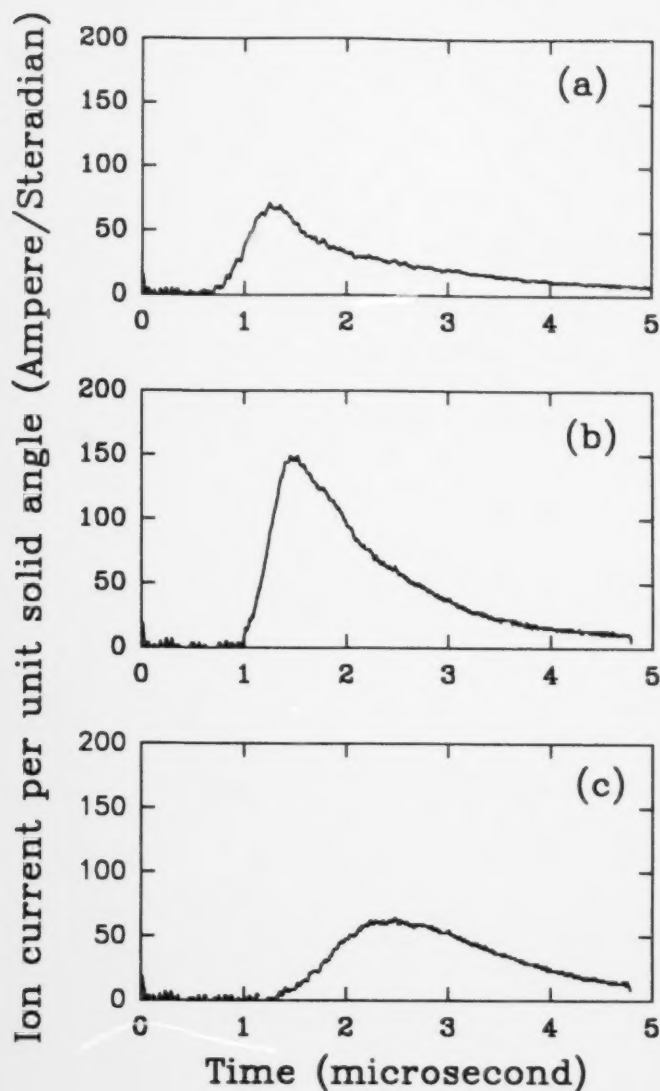


Figure 4.5 Ion current signals observed by a Faraday cup at 25 cm from the target surface and 56° to the target normal. The C plasmas were produced by laser pulses of $0.248 \mu\text{m}$ and 2 ns at irradiance of (a) $I_L = 9.1 \times 10^{12} \text{ W/cm}^2$ (best-focus); (b) $I_L = 2.7 \times 10^{12} \text{ W/cm}^2$ (2 mm-off-focus); and (c) $I_L = 8.7 \times 10^{11} \text{ W/cm}^2$ (4 mm-off-focus).

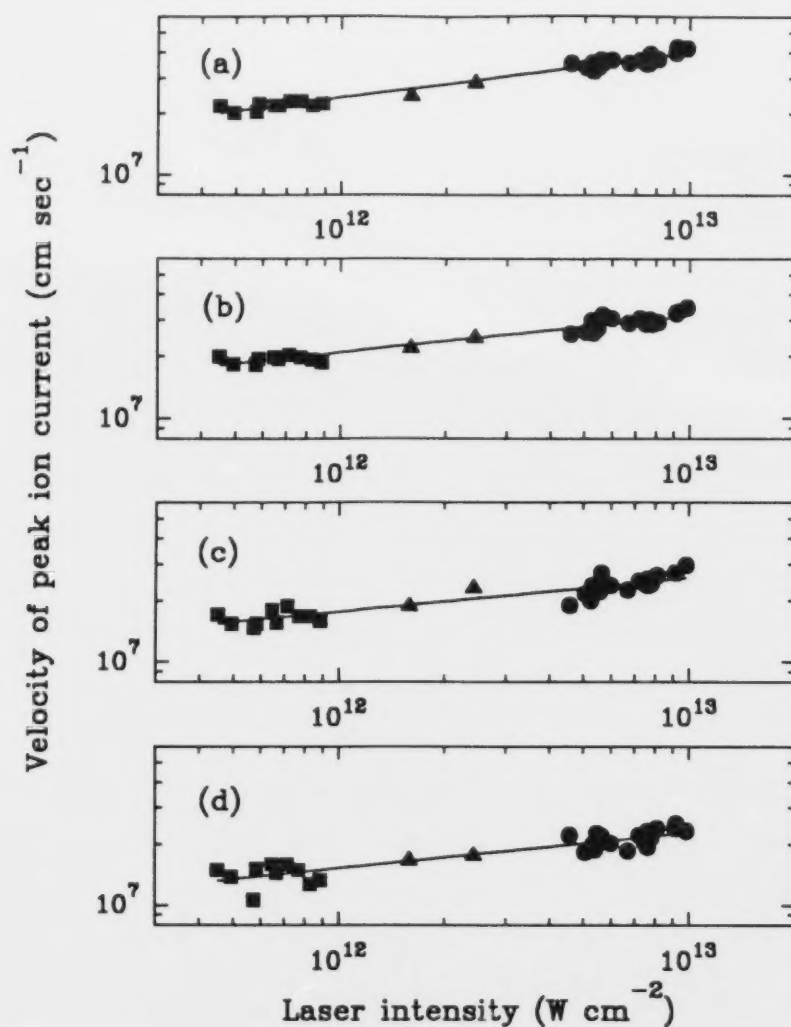


Figure 4.6 Velocity of peak ion current as a function of laser intensity observed at 25 cm from the target surface and (a) 14°, (b) 28°, (c) 42° and (d) 56° to the target normal. The C plasmas were produced by laser pulses of 0.248 μm and 2 ns with three different focal conditions: best focus, 2 mm off focus, and 4 mm off focus (represented by circle, triangle and square respectively).

and B are given in Table 4.1.

Table 4.1 Scaling law of v_{pk} as a function of I_{13}

θ	A	B
14°	3.99 ± 0.47	0.22 ± 0.01
28°	3.16 ± 0.38	0.18 ± 0.01
42°	2.58 ± 0.34	0.17 ± 0.01
56°	2.28 ± 0.31	0.18 ± 0.02

The velocity of peak ion current, v_{pk} , as a function of observation angle for laser intensities of 10^{11} to 10^{13} W/cm² is given in figure 4.7.

4.1.b Charge State Measurements for Short Pulse KrF Laser Produced Carbon Plasma

The ion current for different charge states measured by an electrostatic charge analyzer positioned at a distance 65 cm from the target surface and in a direction $\theta = 20^\circ$ to the target normal is shown in figure 4.8. Results are obtained from data of six shots which have similar pulse length and energy. The average intensity of these shots is 5×10^{12} W/cm². The dominant charge state species are C^{3+} and C^{4+} .

The ion current for different charge states and angles of observation produced by a single laser shot at a nominal laser intensity of 8×10^{12} W/cm² was measured by the electrodynamic charge analyzer positioned at a distance 80 cm from the target surface, as shown in figure 4.9 to 4.11. Each different observation angle is obtained by tilting the target. The dominant charge state species are C^{3+} to C^{5+} for $\theta = 0^\circ$ and 20° . The dominant charge state species is

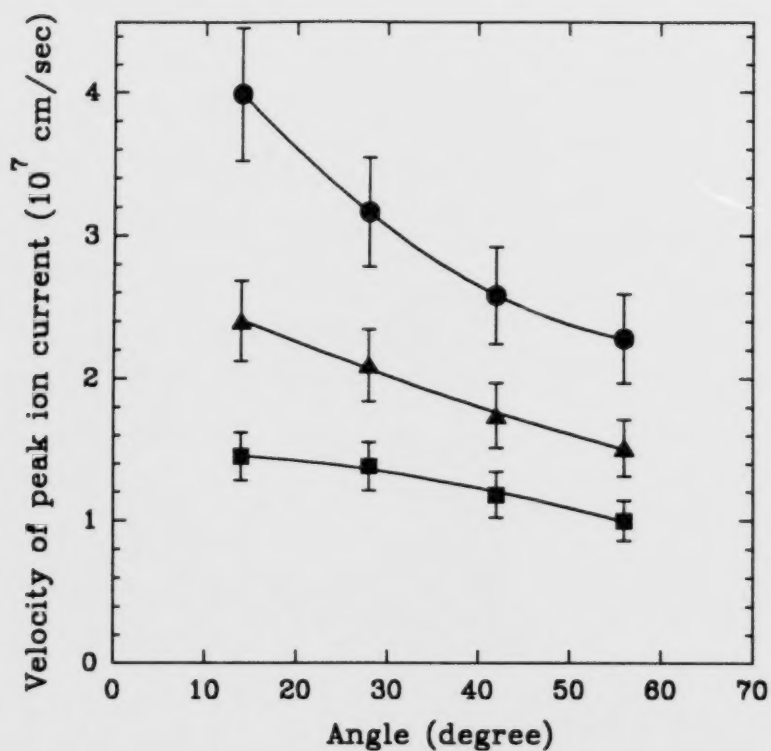


Figure 4.7 Velocity of peak ion current as a function of observation angle at 25 cm from the target surface. The C plasmas were produced by laser pulses of $0.248 \mu\text{m}$ and 2ns at laser irradiances of $1 \times 10^{13} \text{ W/cm}^2$ (circle), $1 \times 10^{12} \text{ W/cm}^2$ (triangle) and $1 \times 10^{11} \text{ W/cm}^2$ (square).

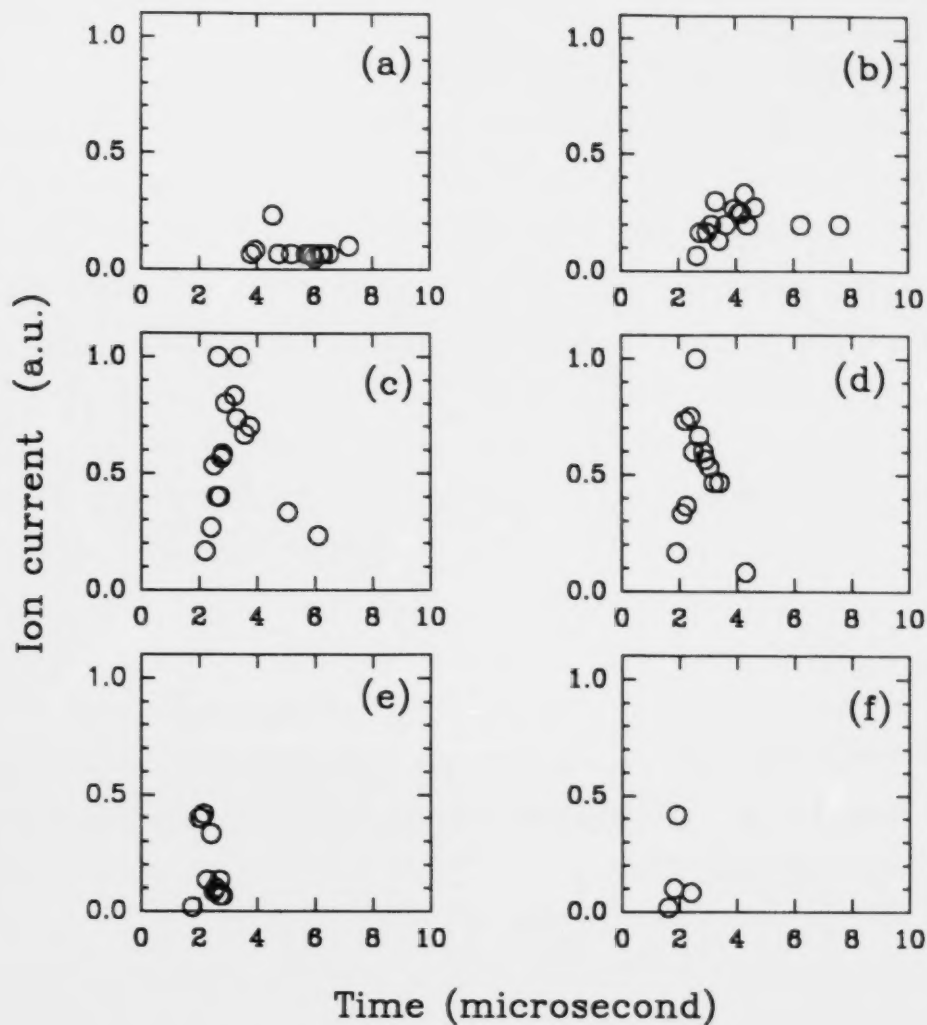


Figure 4.8 Ion current of C^{1+} to C^{6+} (a) to (f) respectively as measured by an electrostatic charge analyzer positioned at a distance 65 cm from the target and in a direction $\theta = 20^\circ$ to the target normal. The $0.248 \mu\text{m}$ laser pulses have a FWHM $\approx 2\text{ns}$. The average laser intensity on target is $5 \times 10^{12} \text{ W/cm}^2$.

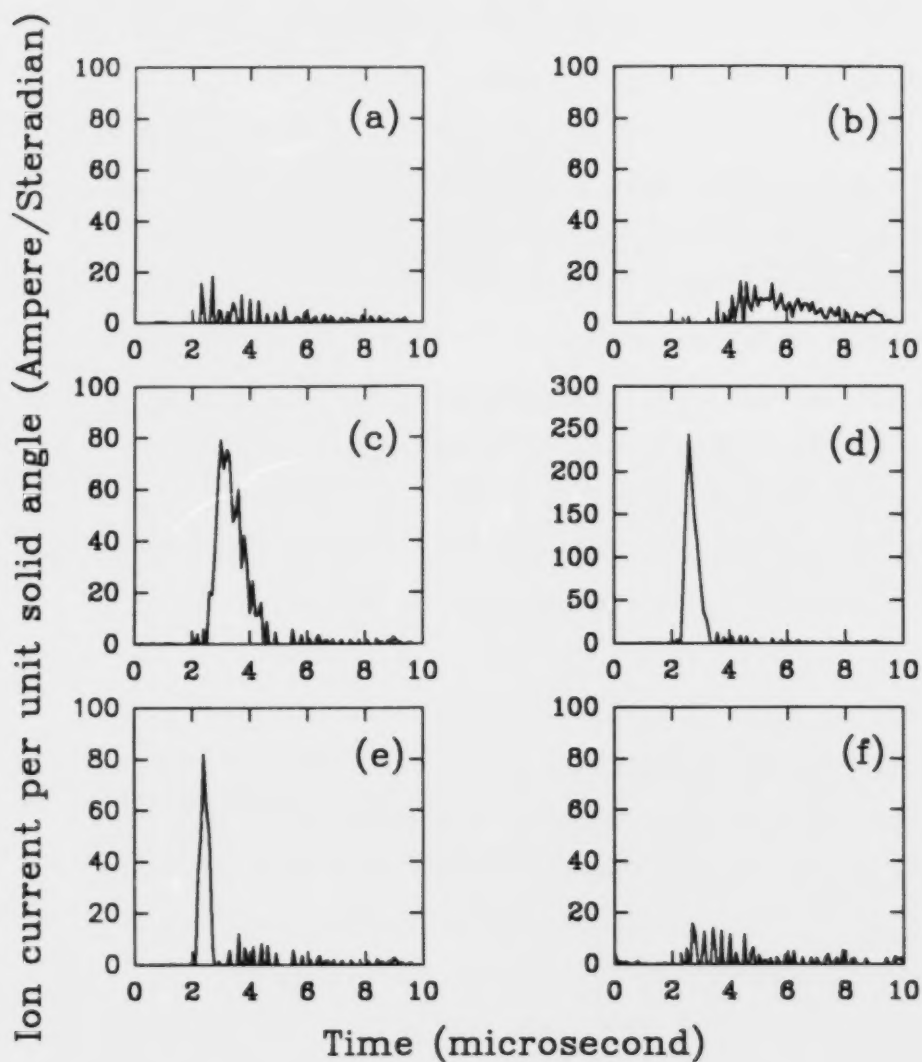


Figure 4.9 Ion current of C^{1+} to C^{6+} , (a) to (f) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 0^\circ$ to the target normal. The $0.248 \mu\text{m}$ laser pulses have a FWHM $\approx 2\text{ns}$. The average laser intensity on target is $8 \times 10^{12} \text{ W/cm}^2$.

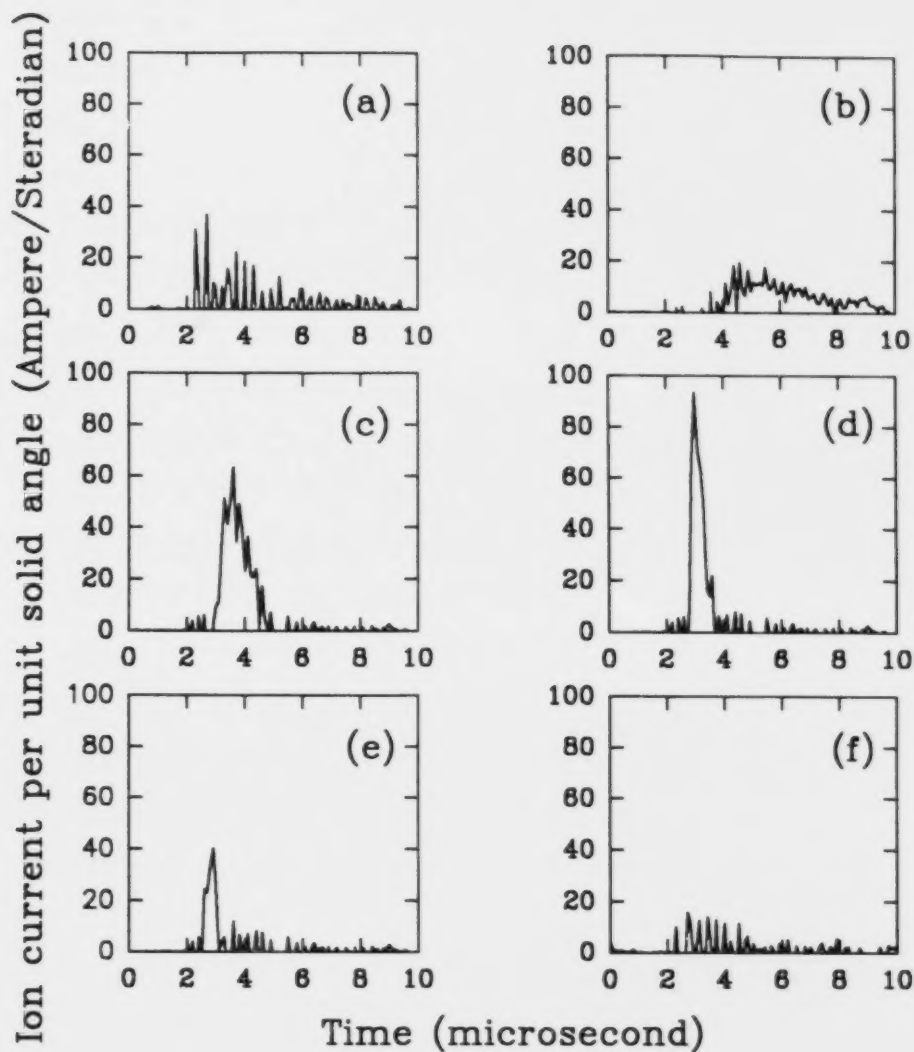


Figure 4.10 Ion current of C^{1+} to C^{6+} , (a) to (f) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 20^\circ$ to the target normal. The $0.248 \mu\text{m}$ laser pulses have a FWHM $\approx 2\text{ns}$. The average laser intensity on target is $8 \times 10^{12} \text{ W/cm}^2$.

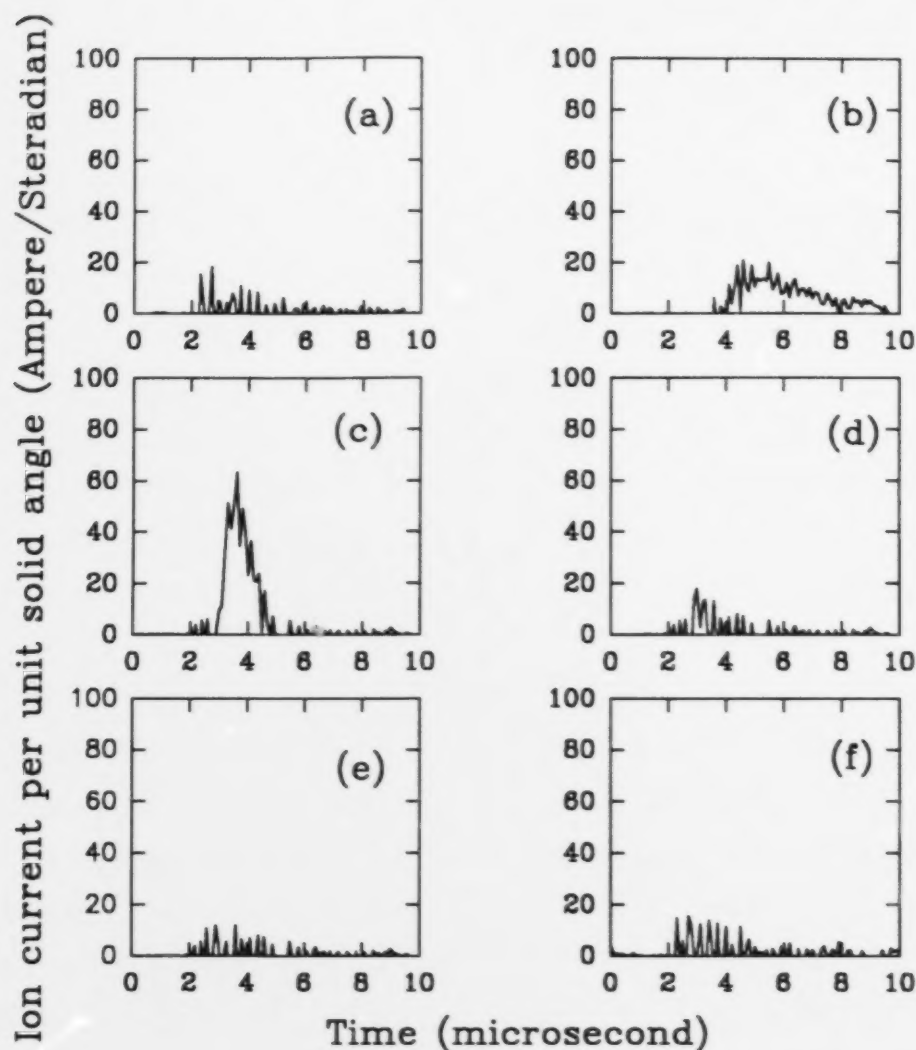


Figure 4.11 Ion current of C^{1+} to C^{6+} , (a) to (f) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 40^\circ$ to the target normal. The $0.248 \mu\text{m}$ laser pulses have a FWHM $\approx 2\text{ns}$. The average laser intensity on target is $8 \times 10^{12} \text{ W/cm}^2$.

C^{3+} for $\theta = 40^\circ$.

The distribution function as a function of velocity for different charge states and angle θ is given in figure 4.12. These results are obtained from data of three shots which have similar pulse length and energy. The average intensities of these shots are $8 \times 10^{12} \text{ W/cm}^2$, $8 \times 10^{12} \text{ W/cm}^2$ and $7 \times 10^{12} \text{ W/cm}^2$ for $\theta = 0^\circ$, 20° and 40° respectively.

The observed variation of average charge state as a function of velocity is approximately linear with no apparent dependence on θ . The best fit for average charge state as a function of velocity is shown in figure 4.13. These results include data from 9 laser shots, 3 for each θ .

4.1.c Ion Emission Measurements for Long Pulse KrF Laser Produced Carbon Plasma

The ion current traces observed in directions of $\theta = 14^\circ$, 28° , 42° and 56° to the target normal at a laser intensity of $1.7 \times 10^{11} \text{ W/cm}^2$ are shown in figure 4.14. As the direction of observation is changed, it is found that the magnitude and velocity of the ion peak decreases with increasing value of θ .

4.1.d Charge State Measurements for Long Pulse KrF Laser Produced Carbon Plasma

The ion current of different charge states measured by an electrostatic charge analyzer positioned at a distance 65 cm from the target surface and in a direction $\theta = 20^\circ$ to the target normal is shown in figure 4.15. These results are obtained from data of six shots which have similar pulse length and energy. The average intensity of these shots is $2 \times 10^{11} \text{ W/cm}^2$. The dominant charge state species are C^{2+} and C^{3+} .

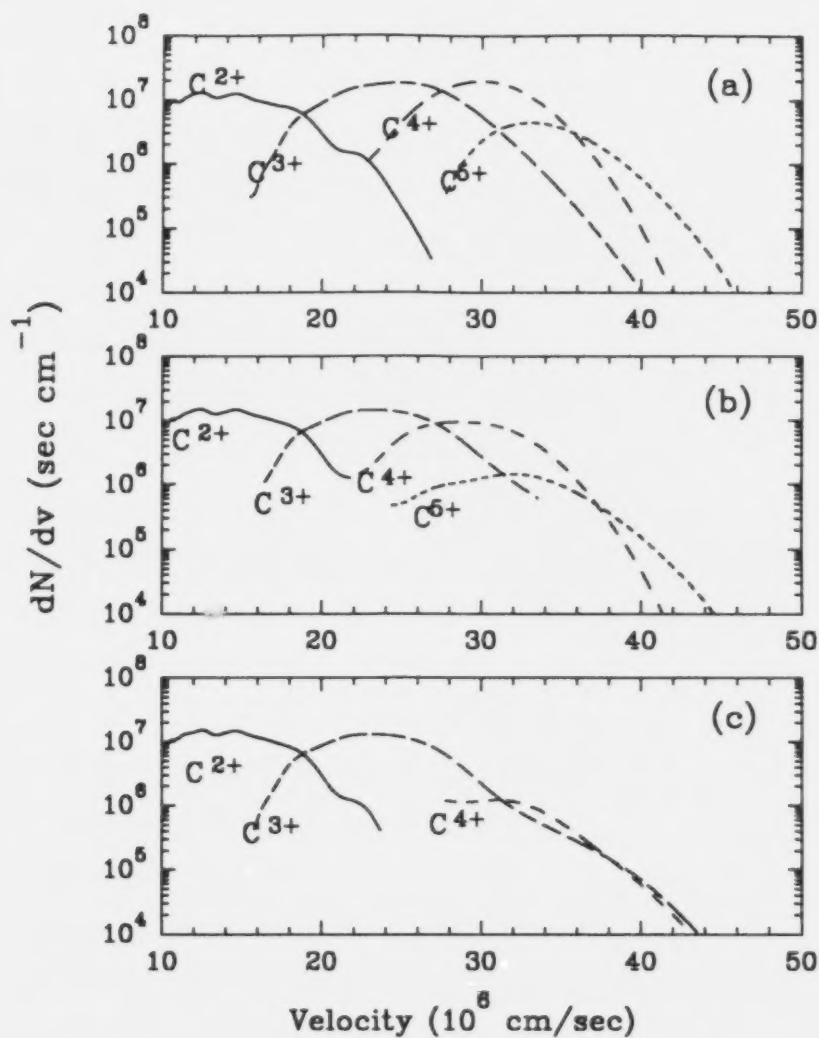


Figure 4.12 Distribution function dN/dv for different charge states and directions (a) $\theta = 0^\circ$ (b) $\theta = 20^\circ$ and (c) $\theta = 40^\circ$ for C plasma as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target. The $0.248 \mu\text{m}$ laser pulses have a FWHM ≈ 2 ns. The average laser intensity on target is $8 \times 10^{12} \text{ W/cm}^2$.

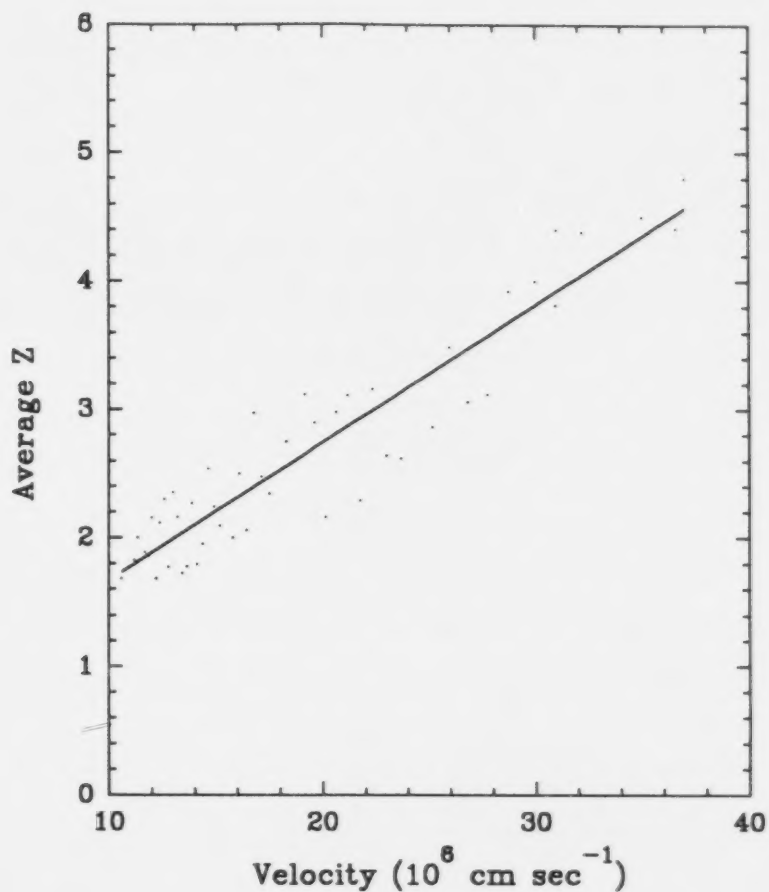


Figure 4.13 Average charge state of carbon plasmas as a function of velocity at a distance 80 cm from the target surface. The plasmas are produced by $0.248 \mu\text{m}$ laser pulses of 2 ns duration. The average laser intensity on target is $8 \times 10^{12} \text{ W/cm}^2$.

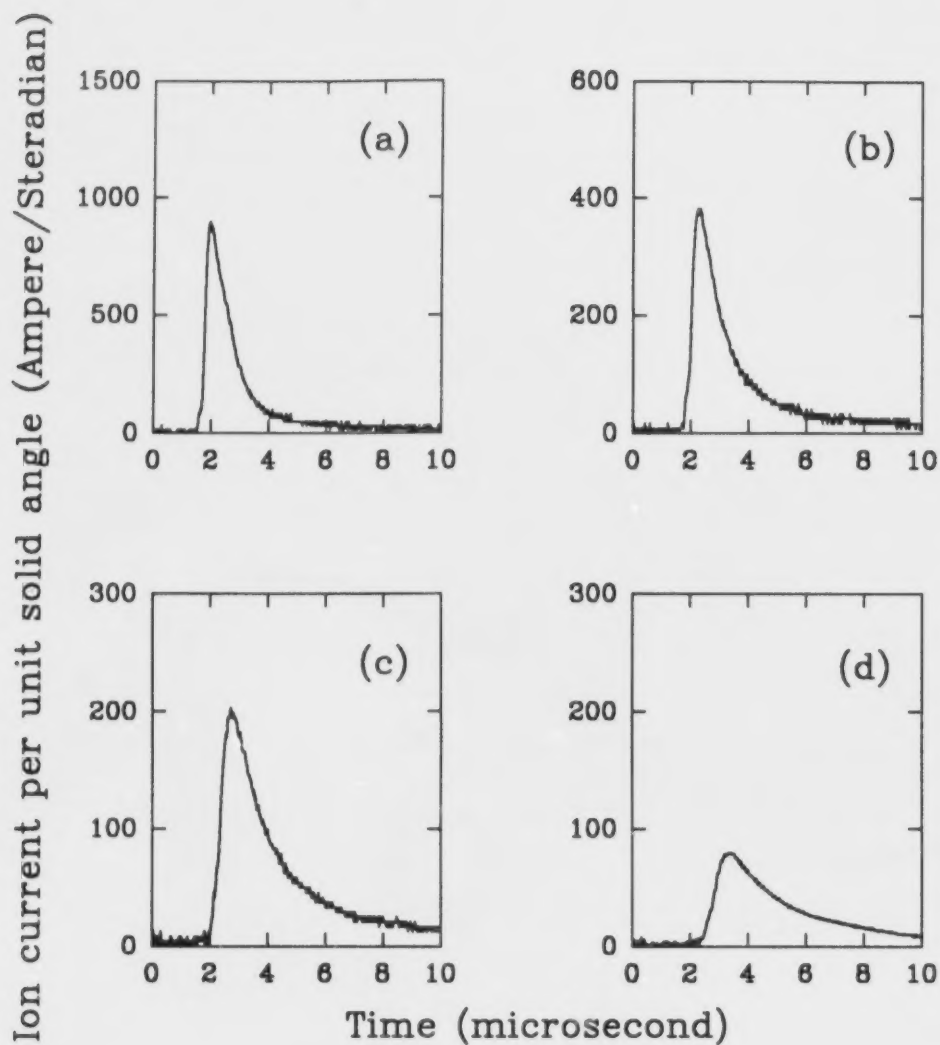


Figure 4.14 Ion current signals observed by a Faraday cup at 25 cm from the target surface for $\theta = 14^\circ$, 28° , 42° and 56° to the target normal, (a) to (d) respectively. The C plasma was produced by a 20 ns $0.248 \mu\text{m}$ laser pulse with an intensity of $1.7 \times 10^{11} \text{W/cm}^2$.

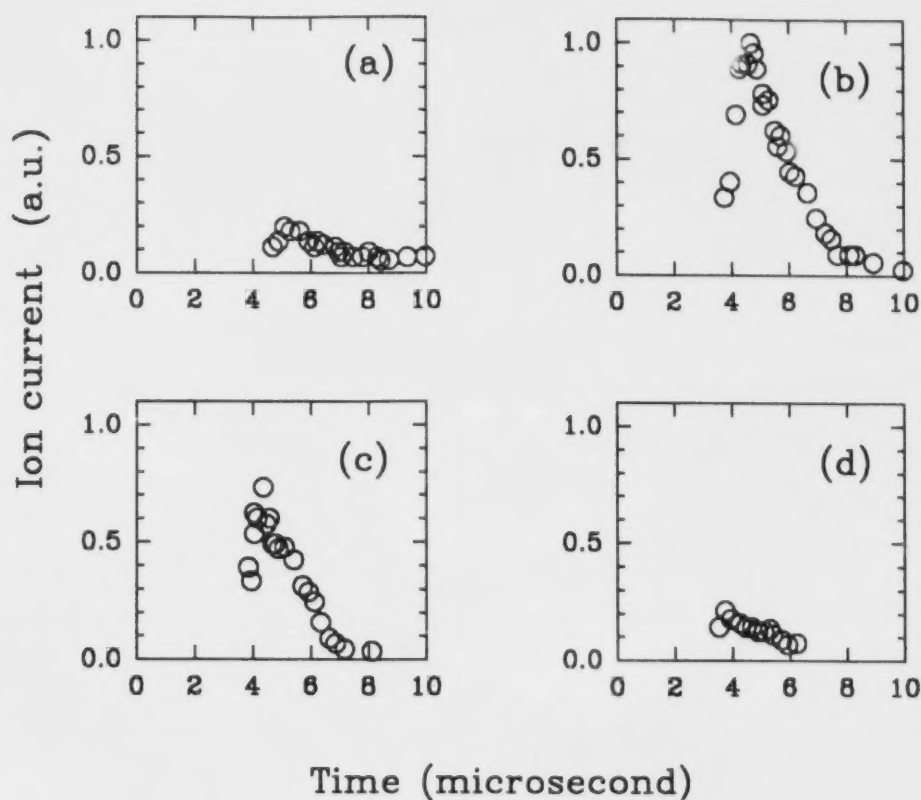


Figure 4.15 Ion current of C^{1+} to C^{4+} , (a) to (d) respectively, as measured by an electrostatic charge analyzer positioned at a distance 65 cm from the target and in a direction $\theta = 20^\circ$ to the target normal. The average intensity on target for the 20 ns 0.248 μm laser pulses is $2 \times 10^{11} \text{ W/cm}^2$.

The ion current of different charge states produced by a single laser shot at a nominal laser intensity of $2 \times 10^{11} \text{ W/cm}^2$ as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target surface at $\theta = 20^\circ$ is shown in figure 4.16. The dominant charge state species are C^{2+} to C^{3+} .

4.1.e Ion Emission Measurements for Ruby Laser Produced Carbon Plasma

The ion current traces observed in directions of $\theta = 0^\circ$, 20° and 45° to the target normal at a nominal laser intensity of $1.0 \times 10^{11} \text{ W/cm}^2$ are shown in figure 4.17. As the direction of observation is changed, it is found that the magnitude and velocity of the ion peak decreases with increasing value of θ .

4.1.f Charge State Measurements for Ruby Laser Produced Carbon Plasma

The ion currents of different charge states measured by an electrostatic charge analyzer positioned at a distance 45 cm from the target surface and in a direction $\theta = 0^\circ$ to the target normal are shown in figure 4.18. These results are obtained from data of six shots which have similar pulse length and energy. The average intensity of these shots is $1 \times 10^{11} \text{ W/cm}^2$. The dominant charge state species are C^{2+} and C^{3+} .

The ion currents of different charge states and angles of observation produced by a single laser shot at a nominal laser intensity of $1 \times 10^{11} \text{ W/cm}^2$, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target surface, shown in figures 4.19 to 4.21. The different observation angles are obtained by tilting the target. The dominant charge state species are C^{2+} to C^{3+} for $\theta = 0^\circ$, 20° , 45° .

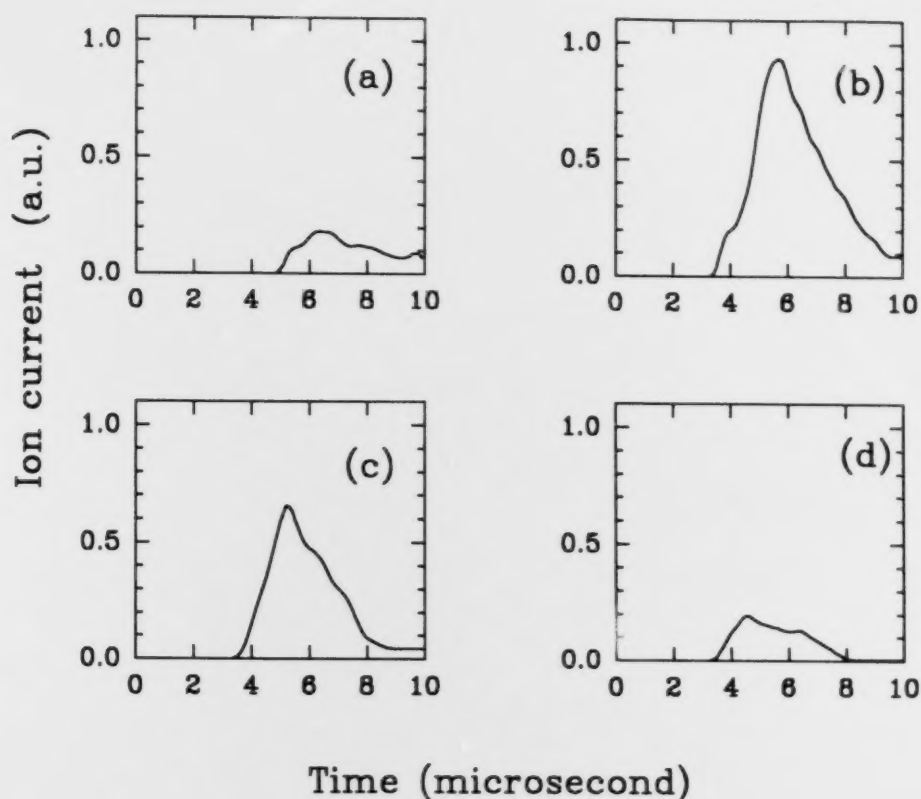


Figure 4.16 Ion current of C^{1+} to C^{4+} , (a) to (d) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 20^\circ$ to the target normal. The intensity on target for the 20 ns 0.248 μm laser pulse is $2 \times 10^{11} \text{ W/cm}^2$.

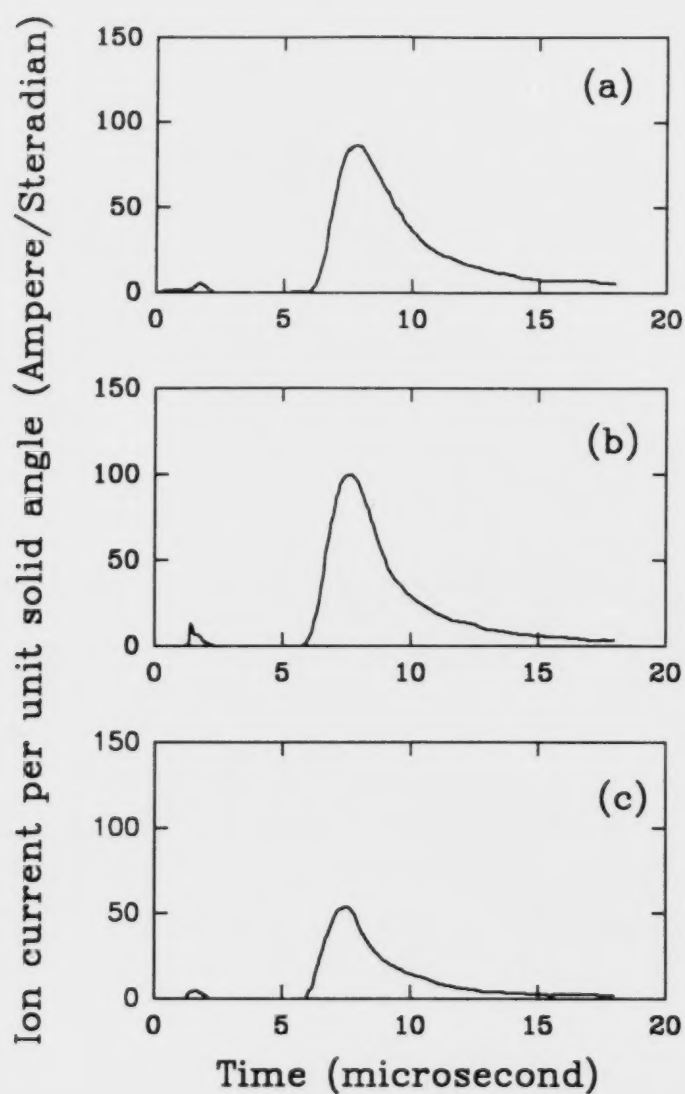


Figure 4.17 Ion current signals observed by Faraday cups at 80 cm from the target surface and $\theta = 0^\circ$ (a), 20° (b) and 45° (c) to the target normal. The C plasma was produced by a 25 ns $0.694 \mu\text{m}$ laser pulse with an intensity of $1 \times 10^{11} \text{W/cm}^2$.

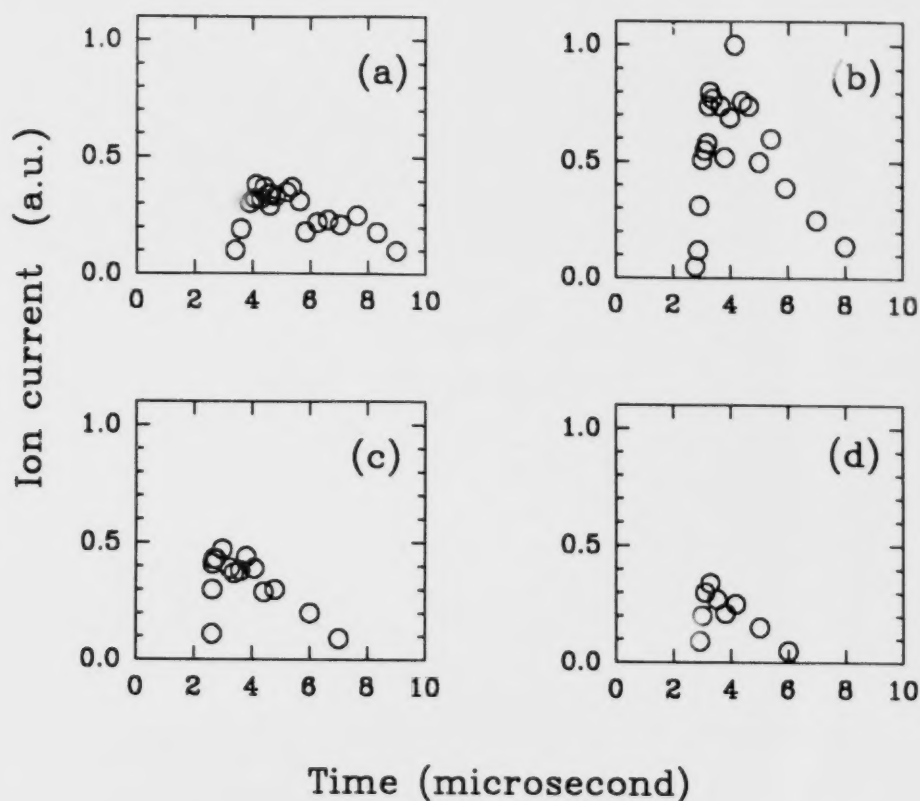


Figure 4.18 Ion current of C^{1+} to C^{4+} , (a) to (d) respectively, as measured by an electrostatic charge analyzer positioned at a distance 45 cm from the target surface and in a direction $\theta = 0^\circ$ to the target normal. The average laser intensity on target for the 25 ns 0.694 μm laser beam is $1 \times 10^{11} \text{ W/cm}^2$.

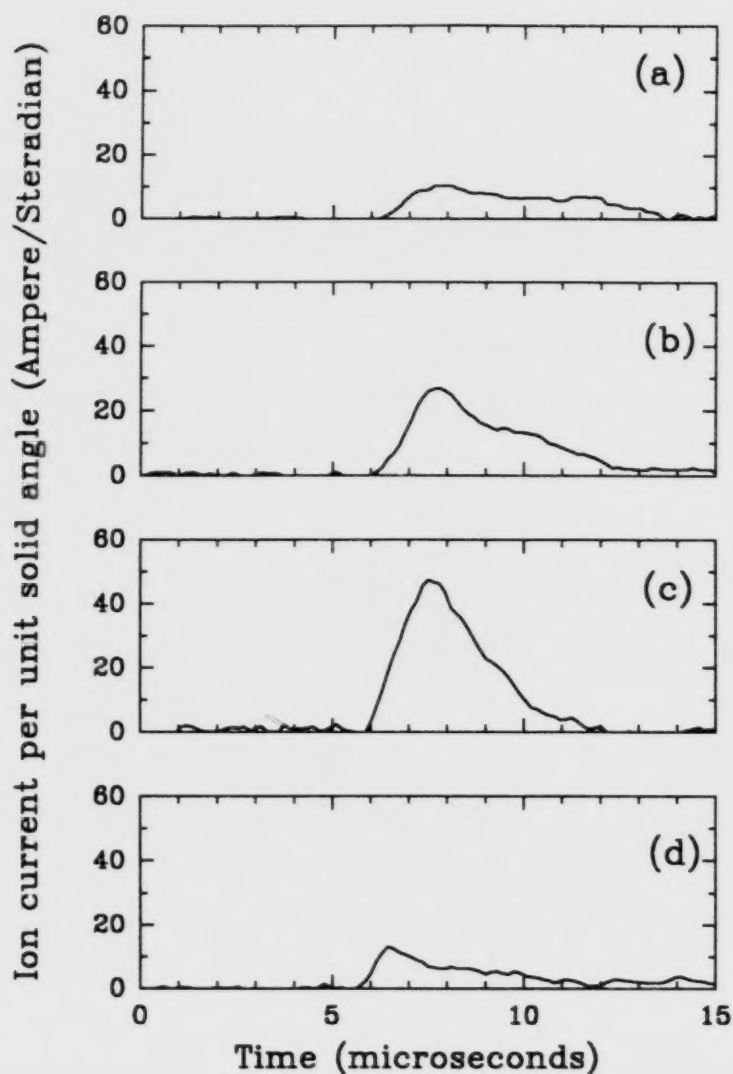


Figure 4.19 Ion current of C^{1+} to C^{4+} , (a) to (d) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 0^\circ$ to the target normal. The laser intensity on target for the 25 ns $0.694 \mu\text{m}$ laser beam is $1 \times 10^{11} \text{ W/cm}^2$.

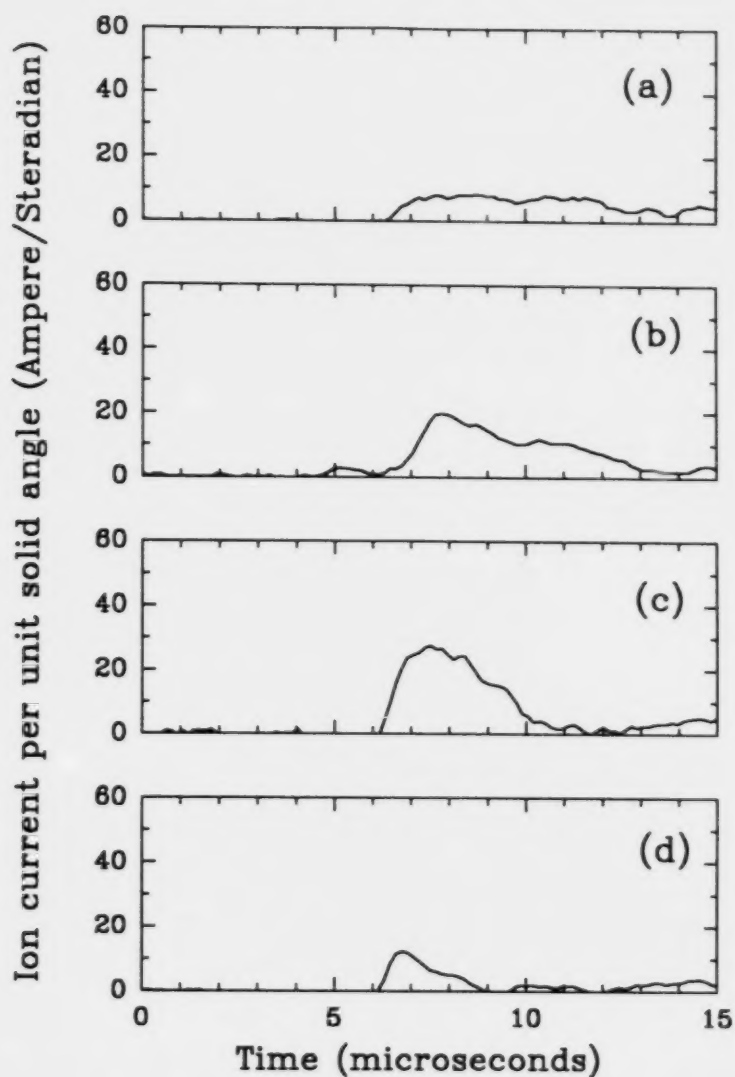


Figure 4.20 Ion current of C^{1+} to C^{4+} , (a) to (d) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 20^\circ$ to the target normal. The laser intensity on target for the 25 ns $0.694 \mu\text{m}$ laser beam is $1 \times 10^{11} \text{ W/cm}^2$.

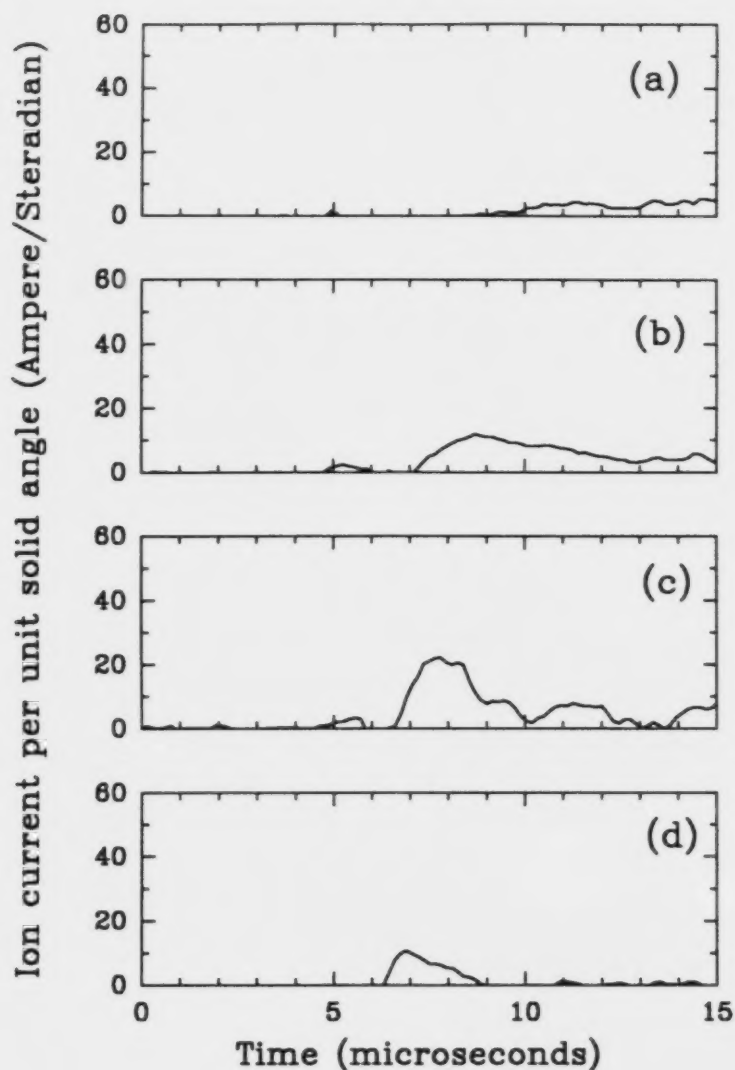


Figure 4.21 Ion current of C^{1+} to C^{4+} , (a) to (d) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 45^\circ$ to the target normal. The laser intensity on target for the 25 ns $0.694 \mu\text{m}$ laser beam is $1 \times 10^{11} \text{W/cm}^2$.

The distribution function as a function of velocity for different charge states at different θ is given in figure 4.22. The results are obtained from data of three shots which have similar pulse length and energy. The average intensity of these shots is $1 \times 10^{11} \text{ W/cm}^2$. The variation of average charge state as a function of velocity is approximately linear with no apparent dependence on θ . The best fit for average charge state as a function of velocity is shown in figure 4.23. The results include data from 9 laser shots, 3 from each θ .

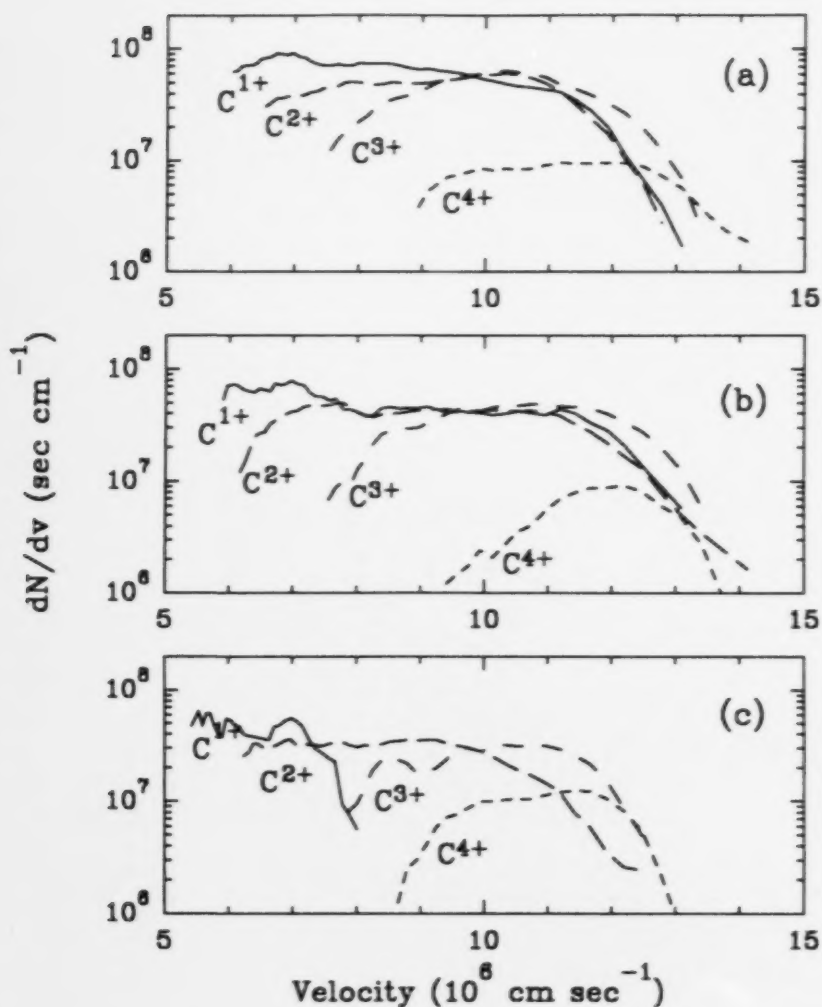


Figure 4.22 Distribution function dN/dv of different charge states and directions (a) $\theta = 0^\circ$, (b) $\theta = 20^\circ$ and (c) $\theta = 45^\circ$ for C plasma as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target. The average laser intensity on target for the 25 ns 0.694 μm laser beam is $1 \times 10^{11} \text{ W/cm}^2$.

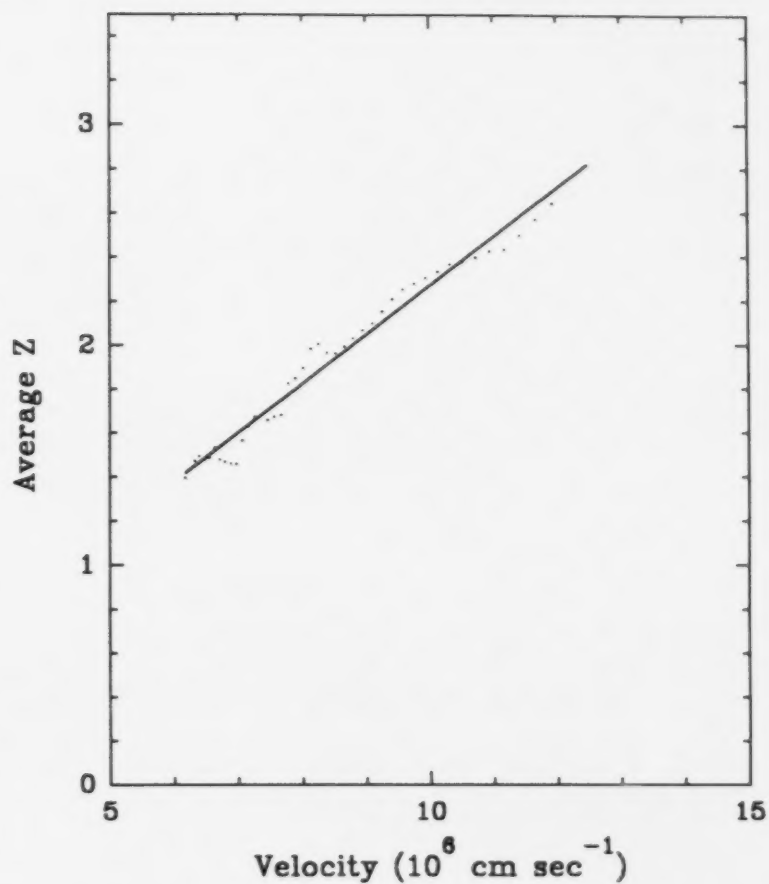


Figure 4.23 Average charge state of C plasma as a function of velocity at a distance 80 cm from the target surface. The plasmas are produced by $0.694 \mu\text{m}$ laser pulses of 25 ns. The average laser intensity on target is $1 \times 10^{11} \text{ W/cm}^2$.

4.2 Ion Expansion Characteristics of Aluminum Plasma

Faraday cup and charge analyzer measurements of ions from an expanding plasma produced by focusing 2 ns, 0.248 and 0.268 μm laser beams on planar aluminum targets in the intensity range of $10^{10} - 10^{13} \text{ W/cm}^2$ are reported in section 4.2.a and 4.2.b; from a 20 ns, 0.248 μm laser beam at intensity of $2 \times 10^{11} \text{ W/cm}^2$ are reported in section 4.2.c and; from a 25 ns, 0.694 μm laser beam at a laser intensity of $1 \times 10^{11} \text{ W/cm}^2$ are reported in section 4.2.d and 4.2.e.

4.2.a Ion Emission Measurements for Short Pulse KrF Laser Produced Aluminum Plasma

The ion current traces observed in directions of $\theta = 14^\circ$, 28° , and 42° to the target normal for different laser intensities are shown in figures 4.24 to 4.26. For $\theta = 14^\circ$, it is observed that in the laser intensity range of $5 \times 10^{11} - 10^{13} \text{ W/cm}^2$ for best-focus (f_0) and 2-mm out-of-focus (f_1), the ion current traces show a two-peak behavior. It is found that the first peak becomes increasingly dominant with increasing laser intensity. As the direction of observation is changed, the second peak becomes relatively more prominent with increasing value of θ . For smaller laser intensity in the large-focal-spot irradiation case (f_2), only one peak was observed with a FWHM value of ($\Delta t/t \approx 0.7 - 1.0$).

In general, the magnitude and velocity of the two ion peaks increase with increasing laser intensity but with a different scaling. As the direction of observation is changed, it is found that the magnitude and velocity of the two ion peaks decrease with increasing value of θ .

The laser intensity scaling of the velocity of the two peaks of the ion current traces observed in directions $\theta = 14^\circ$, 28° , 42° and 67° to the target normal are shown in

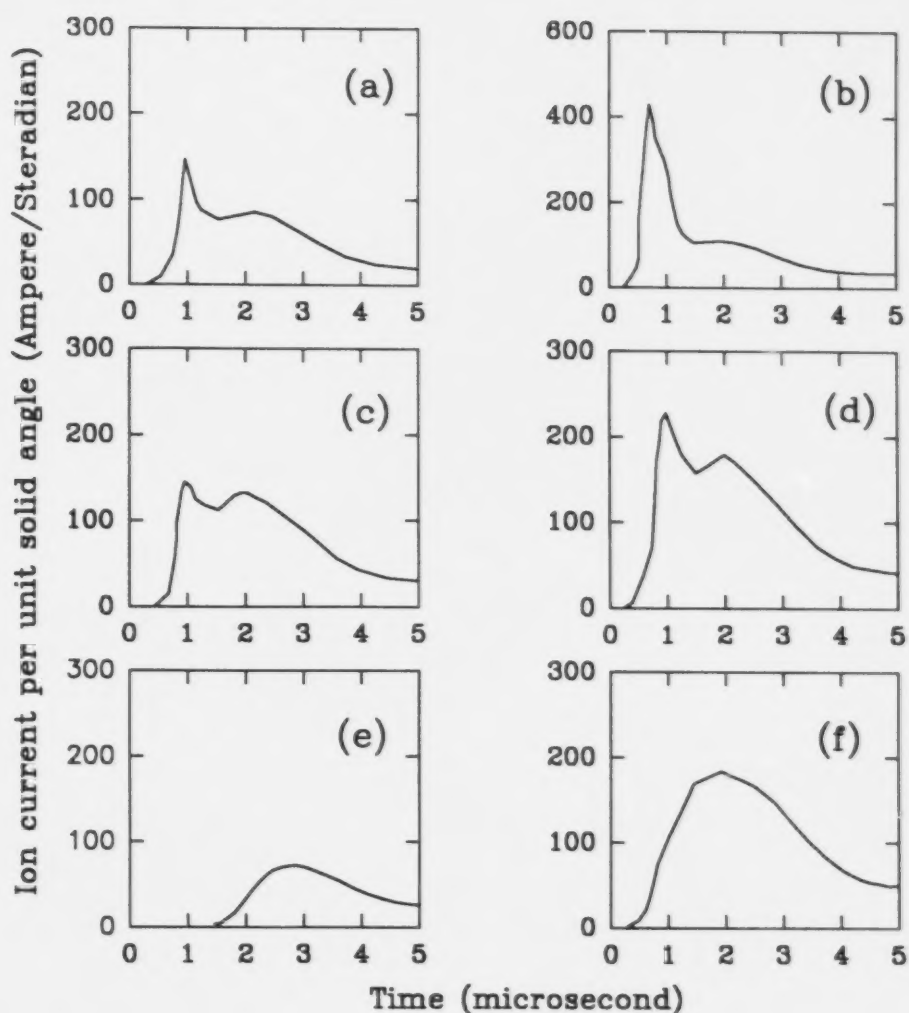


Figure 4.24 Ion current traces for Al plasma observed in a direction of 14° with respect to target normal. The laser intensities (2 ns, $0.268 \mu\text{m}$) for (a)-(f) are 2.7×10^{12} , 7.1×10^{12} , 6.7×10^{11} , 1.1×10^{12} , 1.2×10^{11} and $3.7 \times 10^{11} \text{ W/cm}^2$, respectively. The 90 % energy containing diameters of the focal spots are $74 \mu\text{m}$ (a) and (b), $174 \mu\text{m}$ (c) and (d) and $340 \mu\text{m}$ (e) and (f).

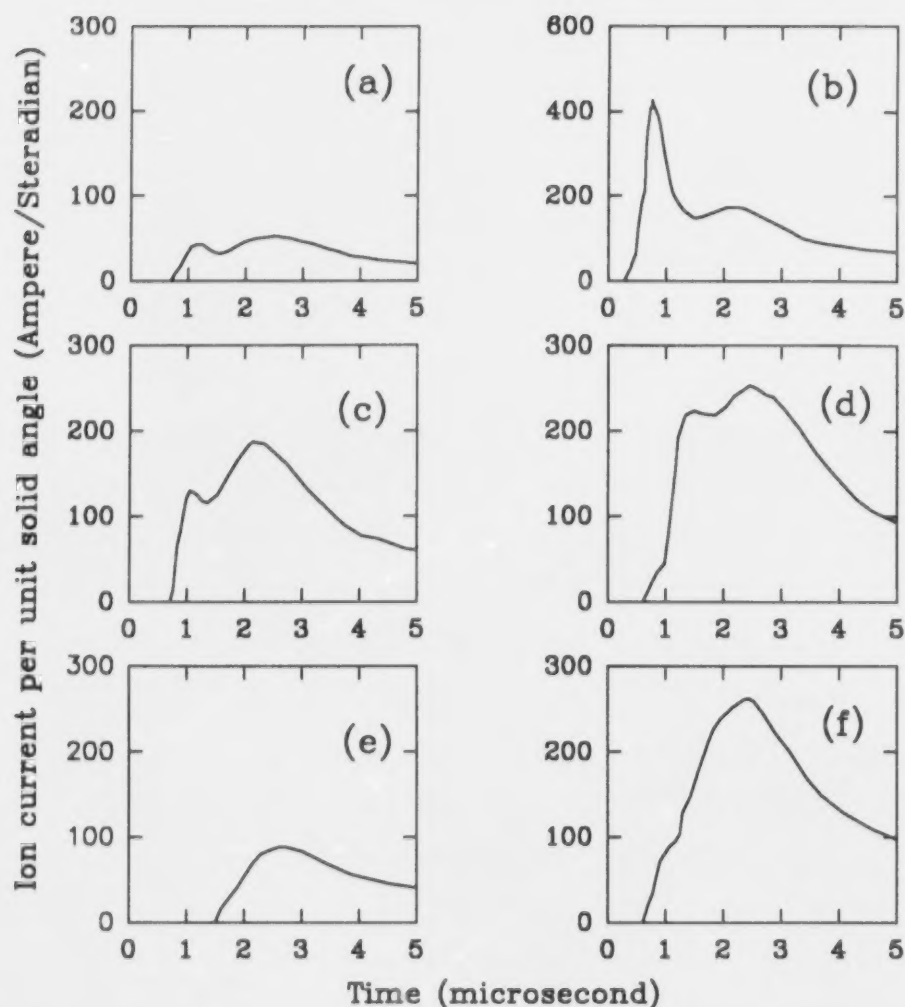


Figure 4.25 Ion current traces for Al plasma observed in a direction of 28° with respect to target normal. The laser intensities (2 ns, $0.268 \mu\text{m}$) for (a)-(f) are 2.7×10^{12} , 7.1×10^{12} , 6.7×10^{11} , 1.1×10^{12} , 1.2×10^{11} and $3.7 \times 10^{11} \text{ W/cm}^2$, respectively. The 90 % energy containing diameters of the focal spots are $74 \mu\text{m}$ (a) and (b), $174 \mu\text{m}$ (c) and (d) and $340 \mu\text{m}$ (e) and (f).

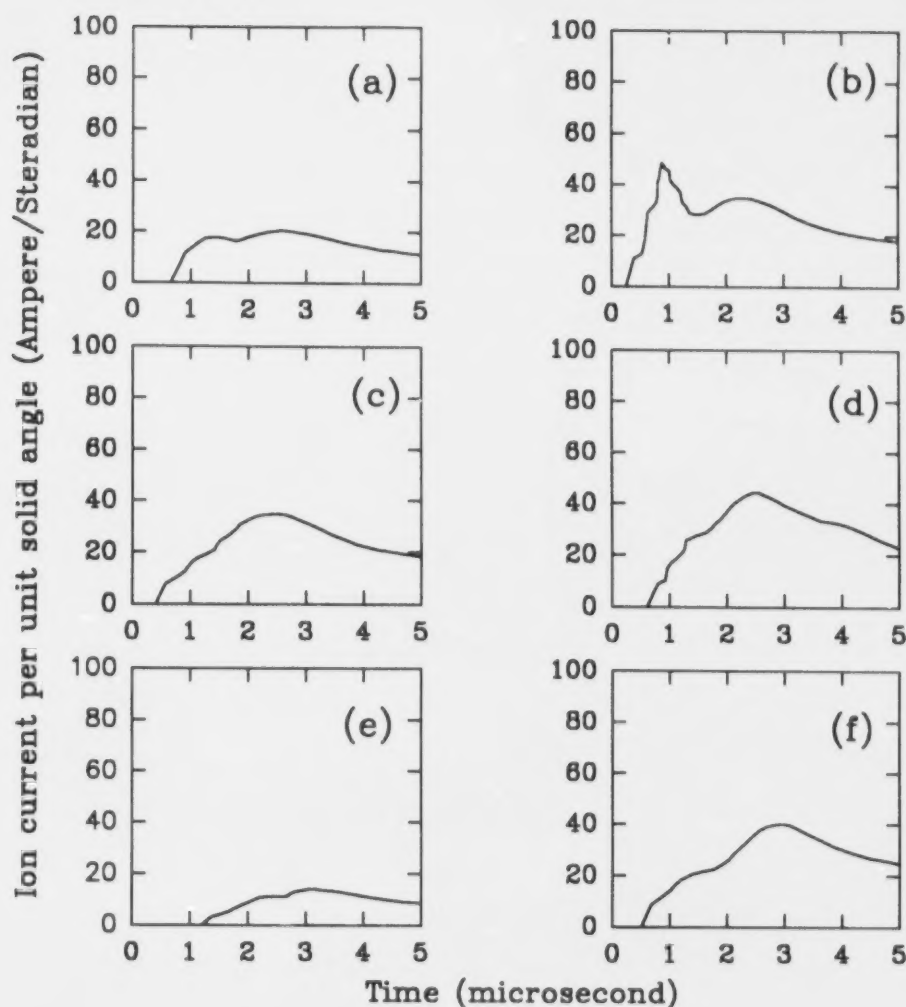


Figure 4.26 Ion current traces of Al plasmas observed in a direction of 42° with respect to target normal. The laser intensities (2 ns, $0.268 \mu\text{m}$) for (a)-(f) are 2.7×10^{12} , 7.1×10^{12} , 6.7×10^{11} , 1.1×10^{12} , 1.2×10^{11} and $3.7 \times 10^{11} \text{ W/cm}^2$, respectively. The 90 % energy containing diameters of the focal spots are $74 \mu\text{m}$ (a) and (b), $174 \mu\text{m}$ (c) and (d) and $340 \mu\text{m}$ (e) and (f).

figure 4.27. The best fit of these data can be summarized in the following equations and table 4.2

$$v_{pk1} = A \times 10^7 I_{13}^B \quad [\text{cm/s}] \quad (4.2)$$

$$v_{pk2} = C \times 10^7 I_{13}^D \quad [\text{cm/s}] \quad (4.3)$$

where v_{pk1} , v_{pk2} and I_{13} are the velocity of the first and second peak of the ion currents and I_{13} is the intensity in unit of 10^{13} W/cm^2 .

Table 4.2 Scaling law of v_{pk1} and v_{pk2} as a function of I_{13}

θ	A	B	C	D
14°	3.57 ± 0.67	0.24 ± 0.01	1.24 ± 0.22	0.04 ± 0.01
28°	3.39 ± 0.64	0.28 ± 0.01	1.11 ± 0.20	0.03 ± 0.01
42°	2.79 ± 0.50	0.29 ± 0.01	1.00 ± 0.18	0.01 ± 0.01
67°	2.06 ± 0.76	0.32 ± 0.04	0.75 ± 0.22	-0.04 ± 0.08

For high intensity, small focal spot irradiation (best focus) the behavior of the ion components contributing to the two peaks were quite different in their characteristics, both in scaling of velocity with laser intensity and in angular distribution of emission. For an observation angle $\theta = 14^\circ$, the scaling for the first peak velocity with laser intensity is $I_L^{0.24}$; the second peak velocity scales much slower - proportional to $I_L^{0.04}$. The former is in good agreement with a scaling of $I_L^{0.22-0.25}$ expected from a steady state ablation model. In terms of angular distribution, ions in the first peak were emitted preferentially closer to normal as compared the ions in the second peak, which showed a more isotropic distribution.

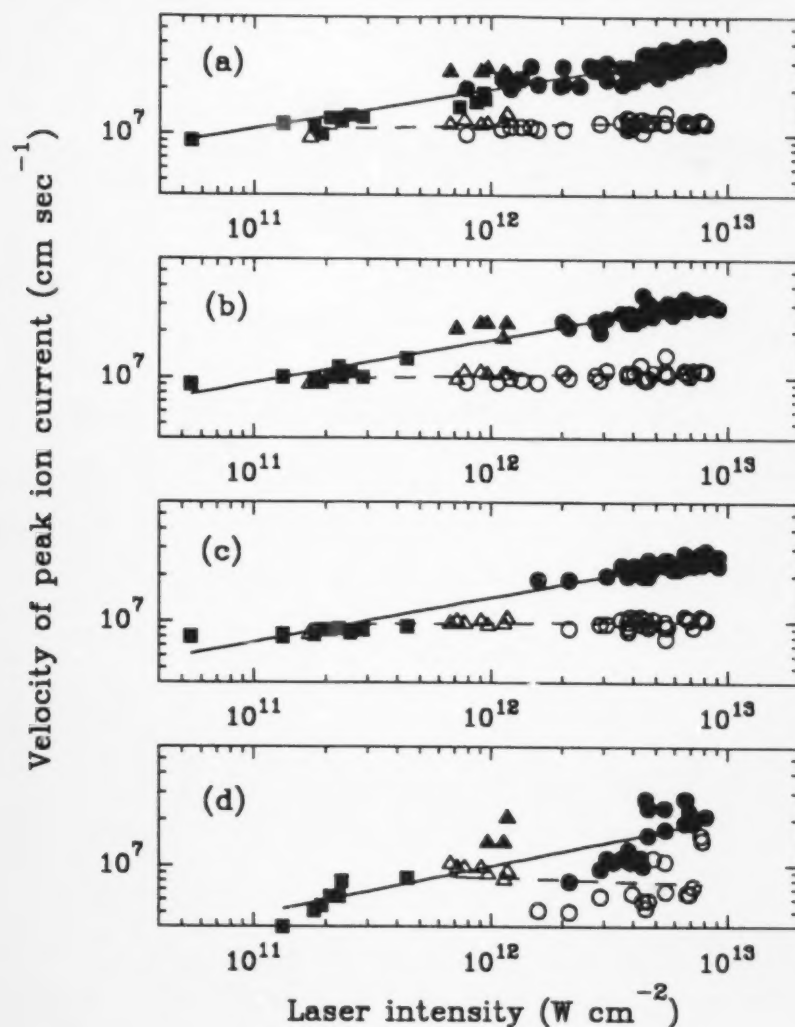


Figure 4.27 Velocity of peak Al ion current as a function of laser intensity observed at (a) $\theta = 14^\circ$, (b) $\theta = 28^\circ$, (c) $\theta = 42^\circ$ and (d) $\theta = 67^\circ$ with respect to target normal. The three different focal spot irradiation conditions for the 2ns and 0.268 μm laser beams are : $d_{90\%} = 74, 174$ and 374 μm (circle, triangle and square symbols respectively.) The velocity for the first and second peaks are represented by solid and hollow symbols respectively.

For large focal spot irradiation (4 mm-off-focus) only one peak is observed and its intensity scaling agrees well with that expected for a steady state ablation model.

The velocity of the two ion peaks v_{pk1} and v_{pk2} as a function of observation angle for laser intensity of 10^{11} to 10^{13} W/cm² is given in figure 4.28.

4.2.b Charge State Measurements for Short Pulse KrF Laser Produced Aluminum Plasma

The ion current of different charge states and angles of observation produced by a single laser shot at a nominal laser intensity of 8×10^{12} W/cm², as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target surface, are shown in figures 4.29 to 4.31. The different observation angles are obtained by tilting the target. The dominant charge state species are Al⁵⁺ to Al⁶⁺ for $\theta = 0^\circ$ and 20° . The dominant charge state species is Al⁴⁺ to Al⁵⁺ for $\theta = 40^\circ$.

The distribution function as a function of velocity for different charge states at different θ is given in figure 4.32. These results are obtained from data of three shots which have similar pulse lengths and energies. The average intensities for these shots are 8×10^{12} W/cm², 8×10^{12} W/cm² and 7×10^{12} W/cm² for $\theta = 0^\circ$, 20° and 40° respectively.

The variation of average charge state as a function of velocity is approximately linear with no apparent dependence on θ . The best fit for average charge state as a function of velocity is shown in figure 4.33. The results include data from 9 laser shots, 3 from each θ .

4.2.c Ion Emission Measurements for Long Pulse KrF Laser Produced Aluminum Plasma

The ion current traces observed in directions $\theta = 14^\circ$,

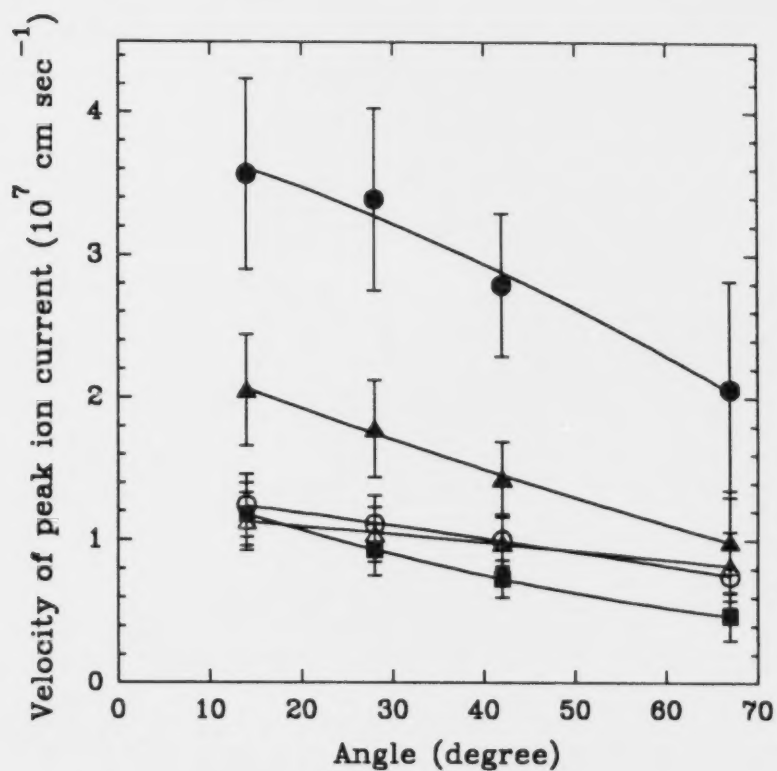


Figure 4.28 Velocity of peak Al ion current as a function of observation angle for (2 ns, 0.268 μm) laser irradiances of 1×10^{13} (circle), 1×10^{12} (triangle) and 1×10^{11} (square) W/cm^2 . The velocities for the first peak and second peak are represented by solid and hollow symbols respectively.

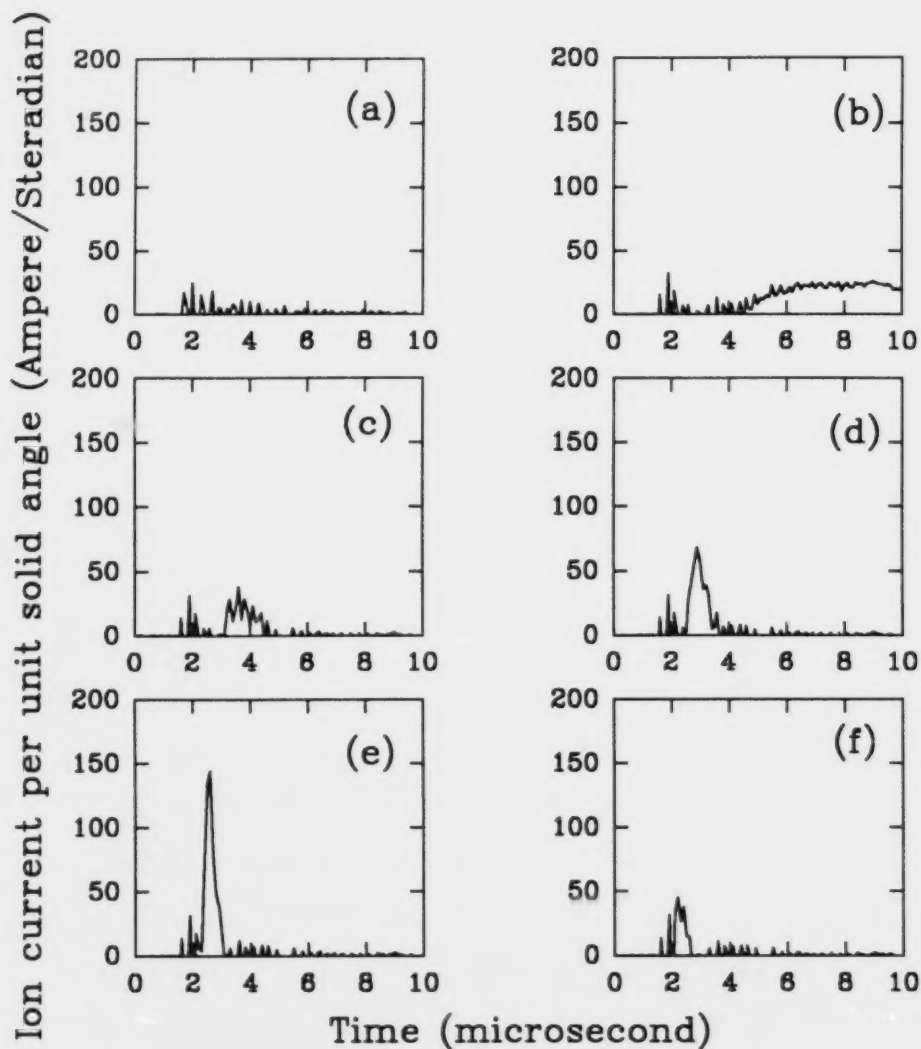


Figure 4.29 Ion current of Al^{3+} to Al^{8+} , (a) to (f) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 0^\circ$ to the target normal. The laser intensity on target for the 2 ns and $0.248 \mu\text{m}$ laser beam is $8 \times 10^{12} \text{ W/cm}^2$.

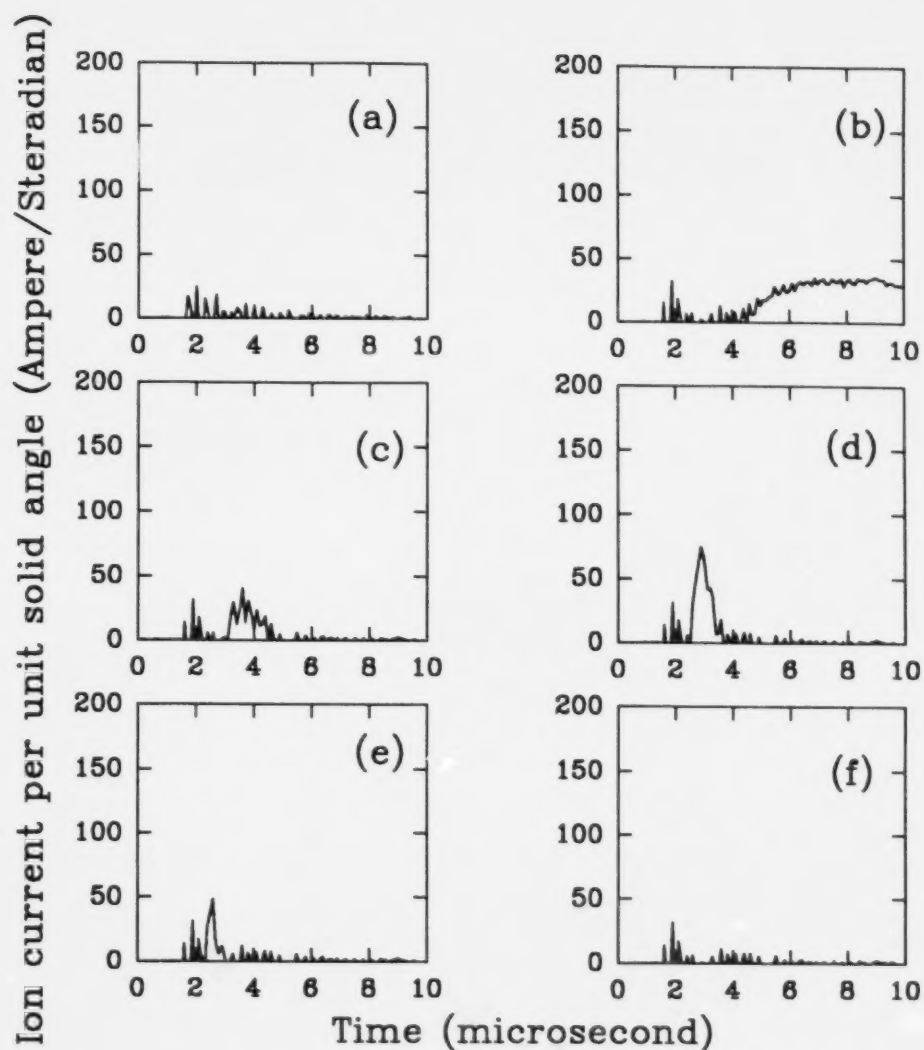


Figure 4.30 Ion current of Al^{3+} to Al^{8+} , (a) to (f) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 20^\circ$ to the target normal. The laser intensity on target for the 2 ns and $0.248 \mu\text{m}$ laser beam is $8 \times 10^{12} \text{ W/cm}^2$.

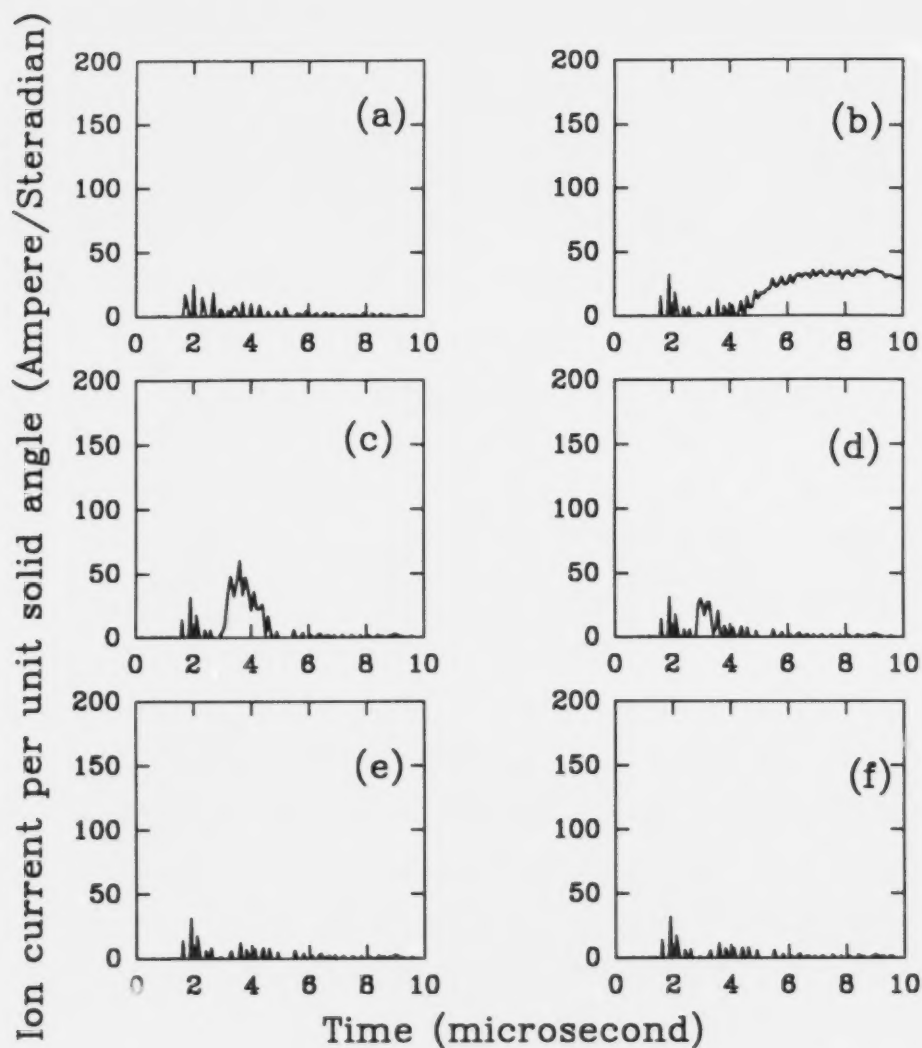


Figure 4.31 Ion current of Al^{3+} to Al^{8+} , (a) to (f) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 40^\circ$ to the target normal. The laser intensity on target for the 2 ns and $0.248 \mu\text{m}$ laser beam is $7 \times 10^{12} \text{ W/cm}^2$.

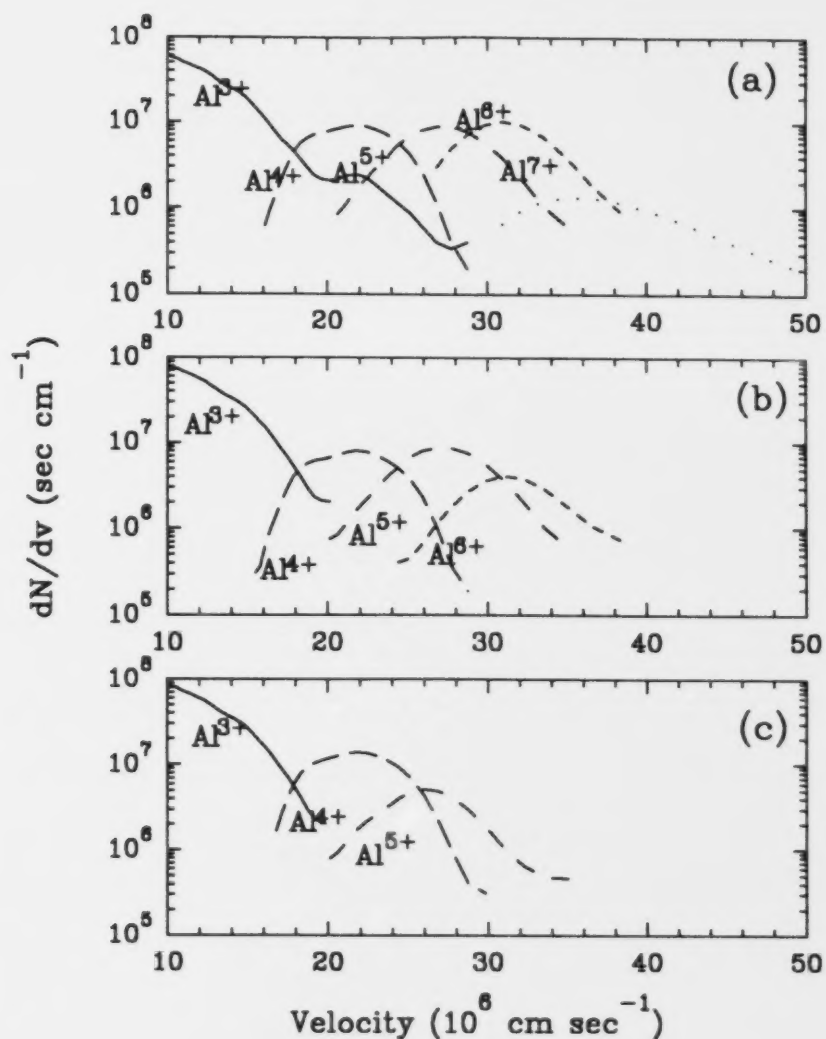


Figure 4.32 Distribution function dN/dv for different charge states and directions of (a) $\theta = 0^\circ$, (b) $\theta = 20^\circ$, and (c) $\theta = 40^\circ$ for Al plasma as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target. The average laser intensity (2 ns, $0.248 \mu\text{m}$) on target is $8 \times 10^{12} \text{ W/cm}^2$.

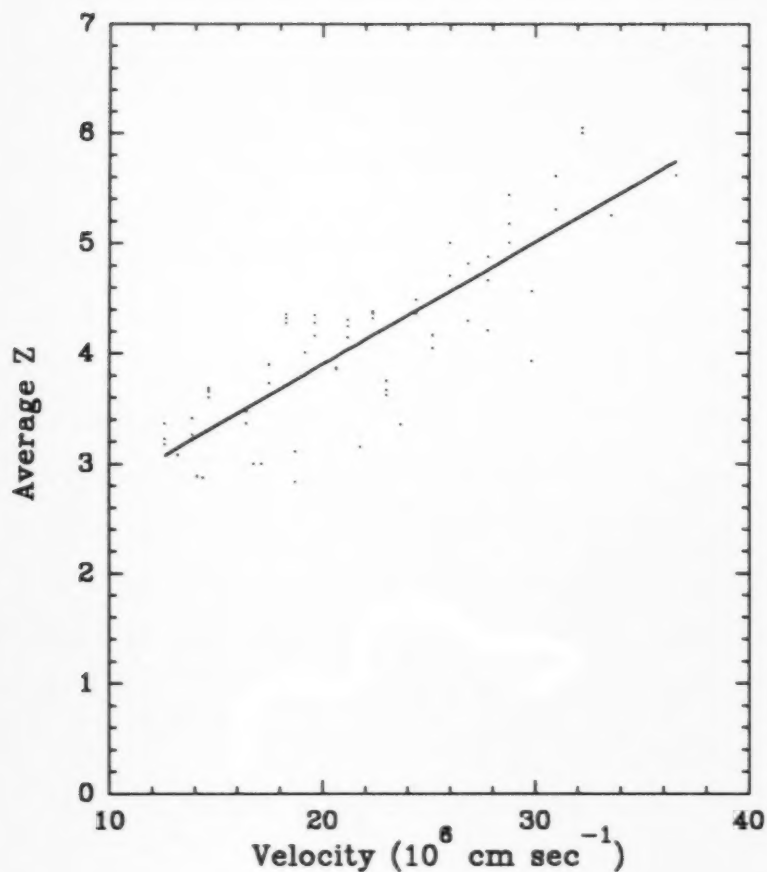


Figure 4.33 Average charge state of Al plasma as a function of velocity at a distance 80 cm from the target surface. The plasmas are produced by $0.248 \mu\text{m}$ laser pulses of 2 ns duration. The average laser intensity on target is $8 \times 10^{12} \text{ W/cm}^2$.

28°, 42° and 56° to the target normal at a laser intensity of $1.7 \times 10^{11} \text{ W/cm}^2$ are shown in figure 4.34. As the direction of observation is changed, it is found the magnitude and velocity of the ion peak decreases with increasing value of θ .

4.2.d Ion Emission Measurements for Ruby Laser Produced Aluminum Plasma

The ion current traces observed in directions $\theta = 0^\circ$, 20° and 45° to the target normal at a nominal laser intensity of $1.0 \times 10^{11} \text{ W/cm}^2$ are shown in figure 4.35. As the direction of observation is changed, it is observed the magnitude and velocity of the ion peak decreased with increasing value of θ .

4.2.e Charge State Measurements for Ruby Laser Produced Aluminum Plasma

The ion current of different charge states and angles of observation produced by a single laser shot at a nominal laser intensity of $1 \times 10^{11} \text{ W/cm}^2$, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target surface, are shown in figures 4.36 to 4.38. The different observation angles are obtained by tilting the target. The dominant charge state species are Al^{2+} to Al^{3+} for $\theta = 0^\circ$, 20° 45° .

The distribution function as a function of velocity for different charge states at different θ is given in figure 4.39. The results are obtained from data of three shots which have similar pulse length and energy. The average intensity of these shots is $1 \times 10^{11} \text{ W/cm}^2$.

The variation of average charge state as a function of velocity is approximately linear with no apparent dependence on θ . The best fit for average charge state as a function of velocity is shown in figure 4.40. The results

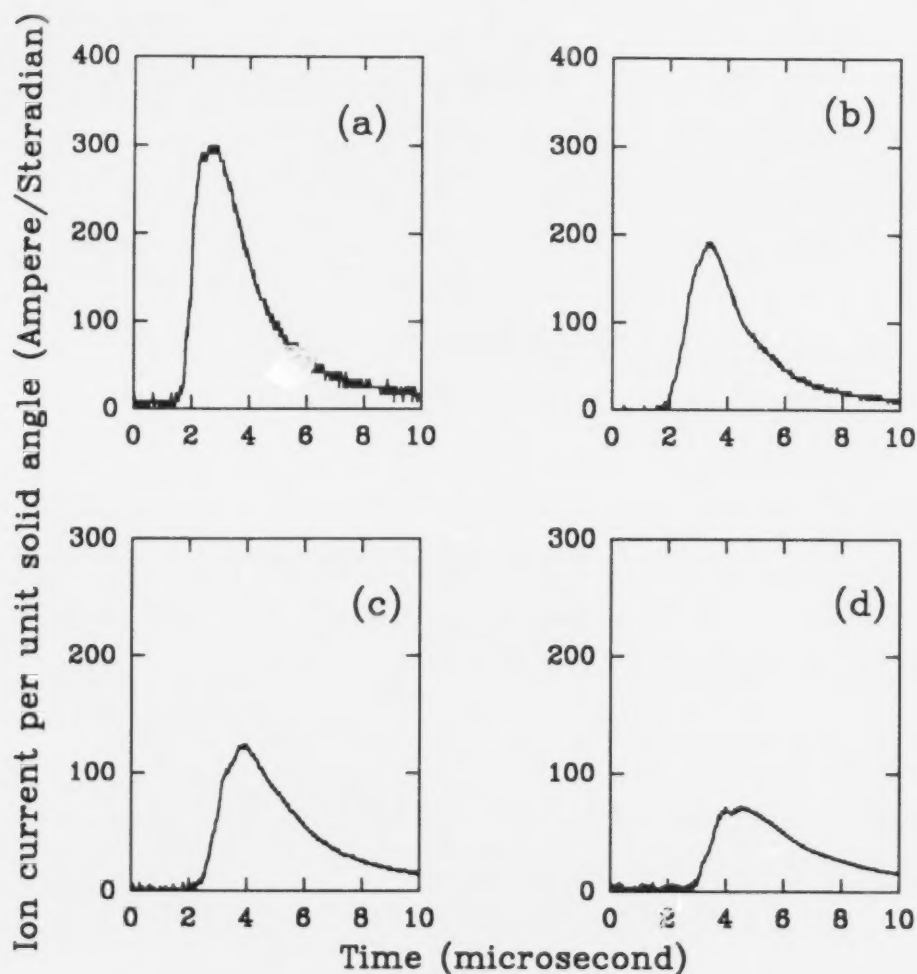


Figure 4.34 Ion current signals observed by Faraday cups at a distance of 25 cm from the target surface at (a) $\theta = 14^\circ$, (b) $\theta = 28^\circ$, (c) $\theta = 42^\circ$ and (d) $\theta = 56^\circ$ to the target normal. The Al plasma was produced by a 20 ns, $0.248 \mu\text{m}$ laser pulse with an intensity of $1.7 \times 10^{11} \text{W/cm}^2$.

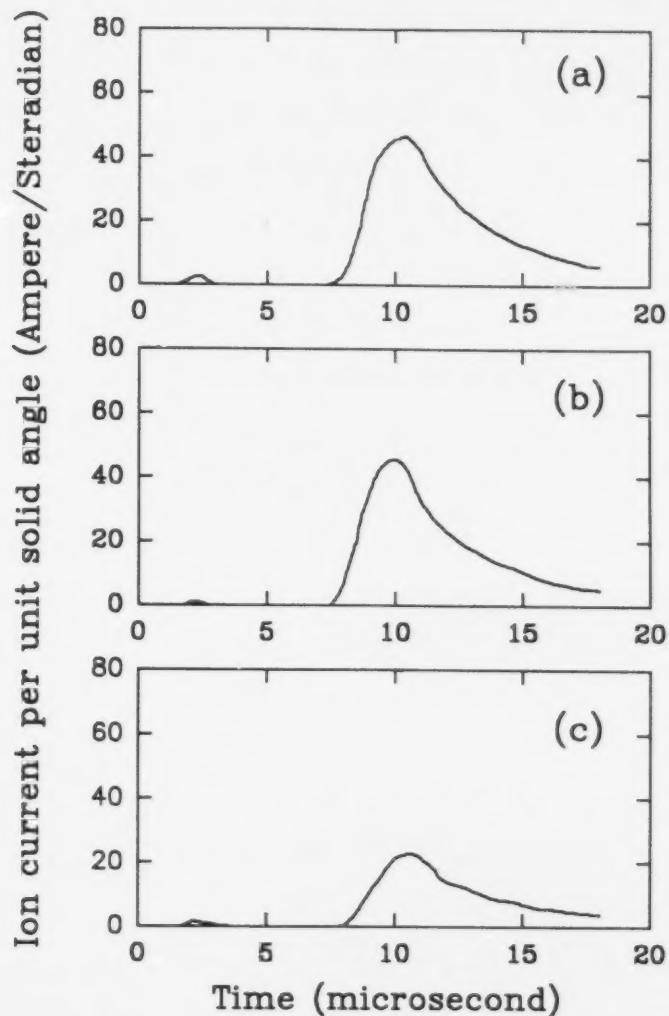


Figure 4.35 Ion current signals observed by Faraday cups at a distance of 80 cm from the target surface at (a) $\theta = 0^\circ$, (b) $\theta = 20^\circ$ and (c) $\theta = 45^\circ$ to the target normal. The Al plasma was produced by a 25 ns, $0.694 \mu\text{m}$ laser pulse with an intensity of $1 \times 10^{11} \text{W/cm}^2$.

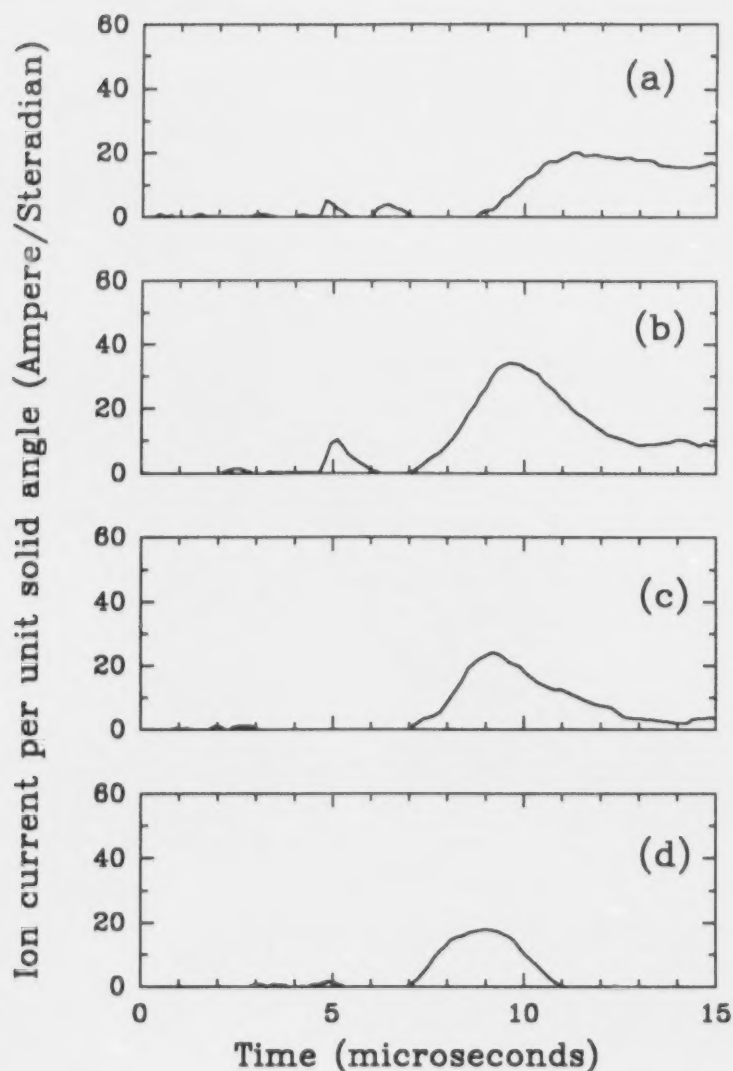


Figure 4.36 Ion current of Al^{1+} to Al^{4+} , (a) to (d) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 0^\circ$ to the target normal. The laser intensity (25 ns, $0.694 \mu\text{m}$) on target is $1 \times 10^{11} \text{ W/cm}^2$.

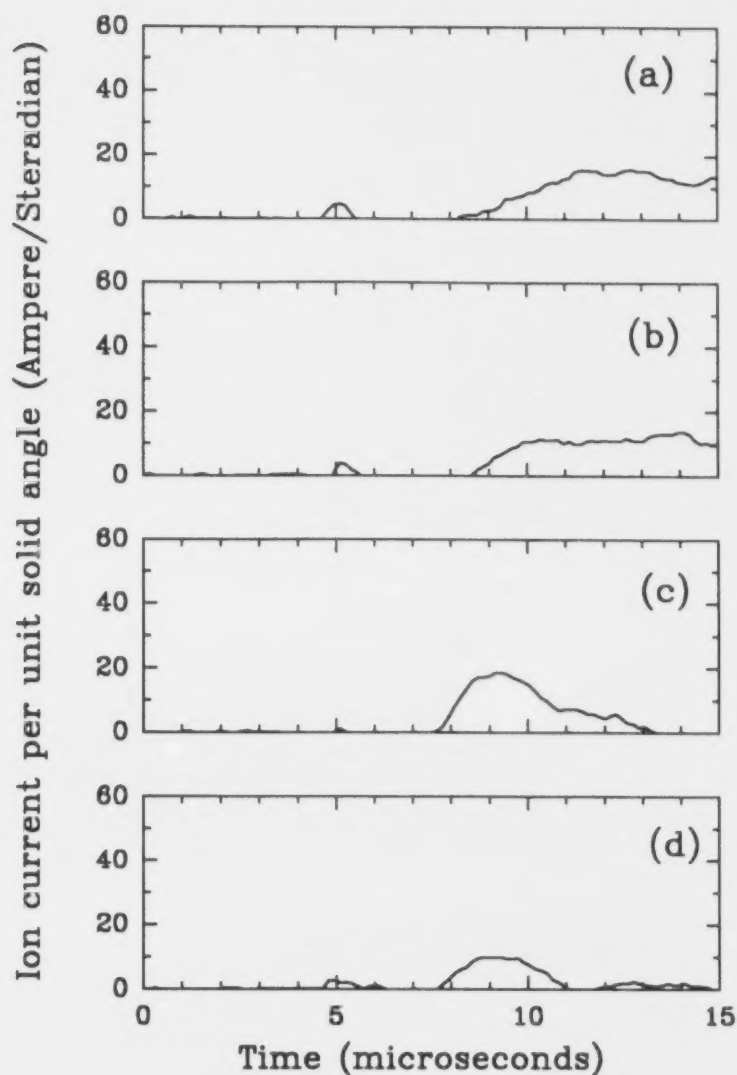


Figure 4.37 Ion current of Al^{1+} to Al^{4+} , (a) to (d) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 20^\circ$ to the target normal. The laser intensity (25 ns, $0.694 \mu\text{m}$) on target is $1 \times 10^{11} \text{ W/cm}^2$.

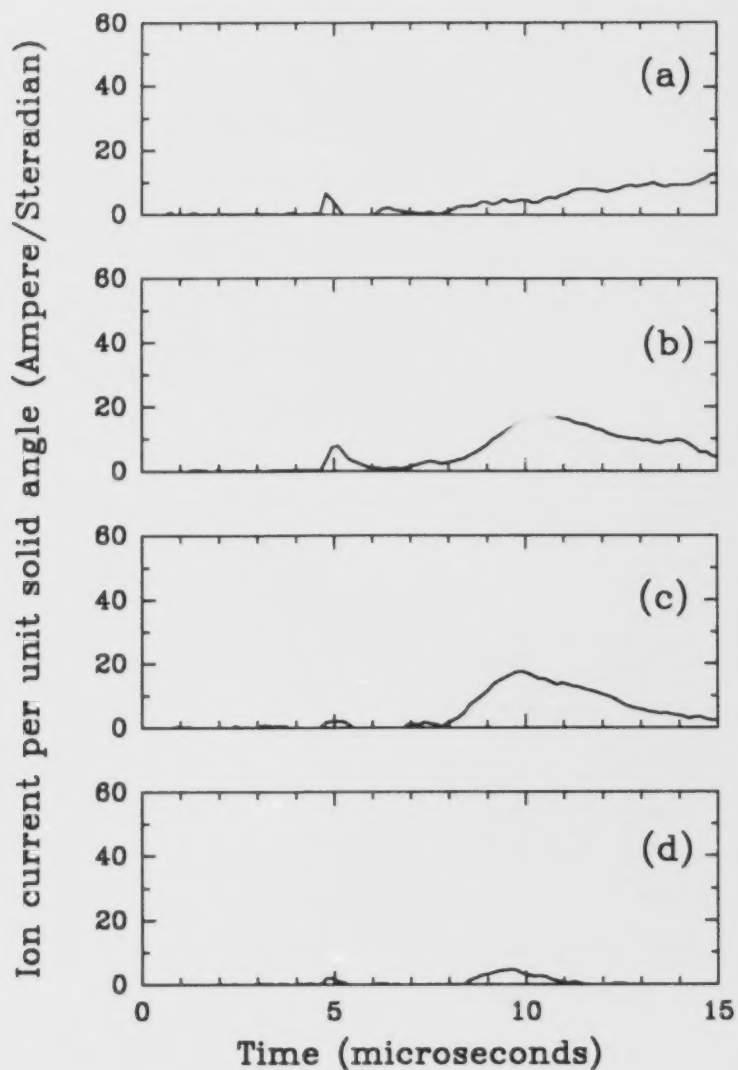


Figure 4.38 Ion current of Al^{1+} to Al^{4+} , (a) to (d) respectively, as measured by an electrodynamic charge analyzer positioned at a distance 80 cm from the target and in a direction $\theta = 45^\circ$ to the target normal. The laser intensity (25 ns, $0.694 \mu\text{m}$) on target is $1 \times 10^{11} \text{ W/cm}^2$.

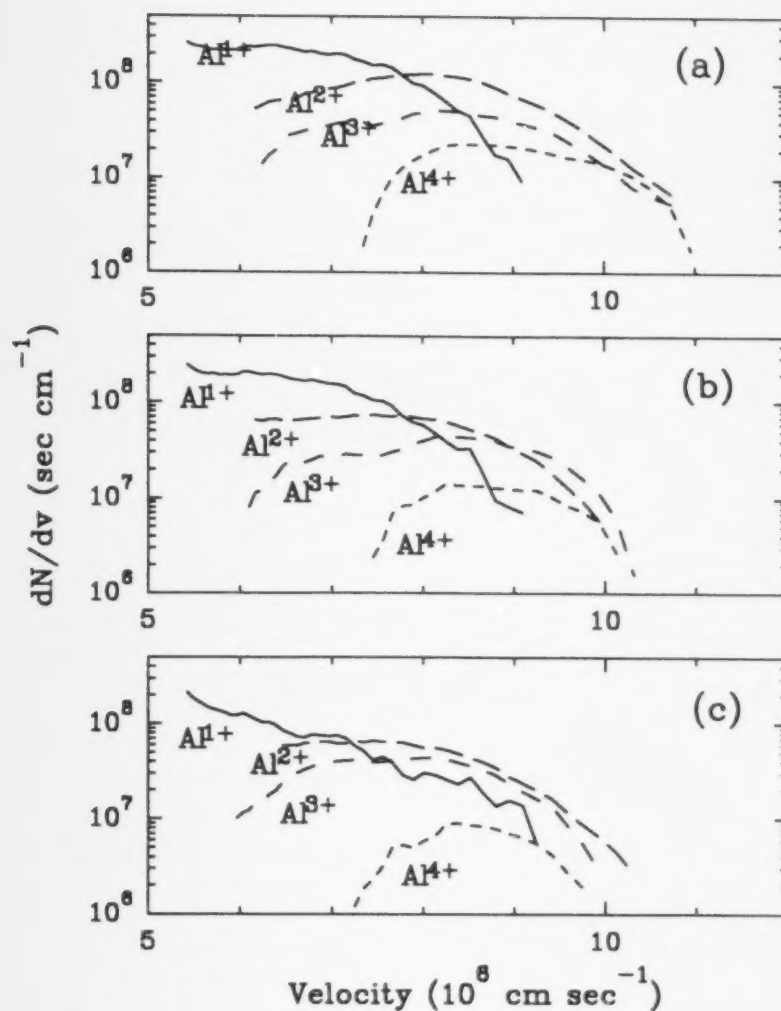


Figure 4.39 Distribution function dN/dv of different charge states and directions of (a) $\theta = 0^\circ$, (b) $\theta = 20^\circ$ and (c) $\theta = 45^\circ$ for an Al plasma as measured by an electrodynamic charge analyzer positioned at 80 cm from the target surface. The average laser intensity (25 ns, $0.694 \mu\text{m}$) on target is $1 \times 10^{11} \text{ W/cm}^2$.

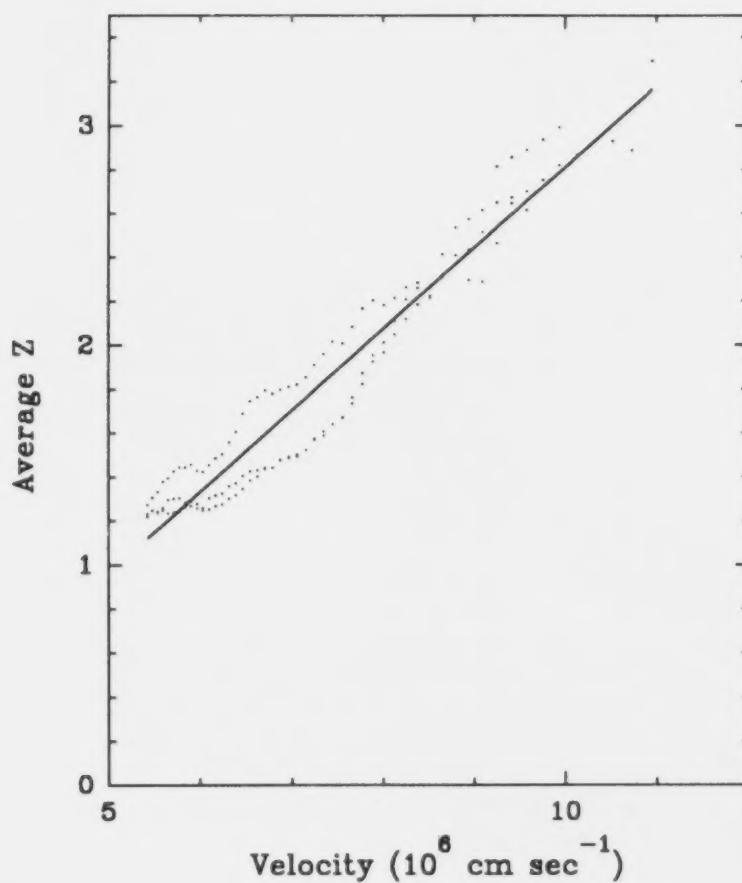


Figure 4.40 Average charge state of Al plasma as a function of velocity at a distance 80 cm from the target surface. The plasmas are produced by $0.694 \mu\text{m}$ laser pulses of 25 ns. The average laser intensity on target is $1 \times 10^{11} \text{ W/cm}^2$.

include data from 9 laser shots, 3 from each θ .

4.3 Ion Expansion Characteristics of Gold Plasma

Faraday cup measurements of ions from an expanding plasma produced by focusing a 2 ns 0.268 μm laser beam onto planar gold targets in the laser intensity range of $10^{10} - 10^{13} \text{ W/cm}^2$ are reported here.

The ion current traces observed in directions $\theta = 14^\circ$, and 28° to the target normal for different laser intensities are shown in figures 4.41 and 4.42. In general, the ion-current signals showed two components which varied with laser intensity and angle of emission.

A number of consistent characteristics are observed in the ion-current signals close to normal incidence ($\theta = 14^\circ$). Firstly, the overall magnitude of ion current decreases as the intensity increases for a fixed laser energy incident on the target. This can be seen from figure 4.41, where the focal spot is varied but the incident laser energy is kept constant at $0.9 \pm 0.1 \text{ J}$. Secondly, the first ion peak is smaller at higher intensities than at lower intensities, again for fixed energy on target. These combined observations indicate that the energy per unit solid angle, which is carried by the ions, is decreasing as the intensity is increasing.

As the angle of observation is increased, the second ion peak decreases more rapidly than the first ion peak when compared to $\theta = 14^\circ$. For $\theta > 42^\circ$ the signal for ion velocity corresponding to the second peak was quite small, so that only the first peak was distinctly seen. These results indicate that ions in the first velocity peak are emitted over a broad angular distribution, while those in the second velocity peak are emitted in a narrow cone of angle of $\approx 30^\circ$ for the best focus irradiation condition. The latter angular spread was smallest for the large-focal-spot (4 mm out of focus) case. This is to be

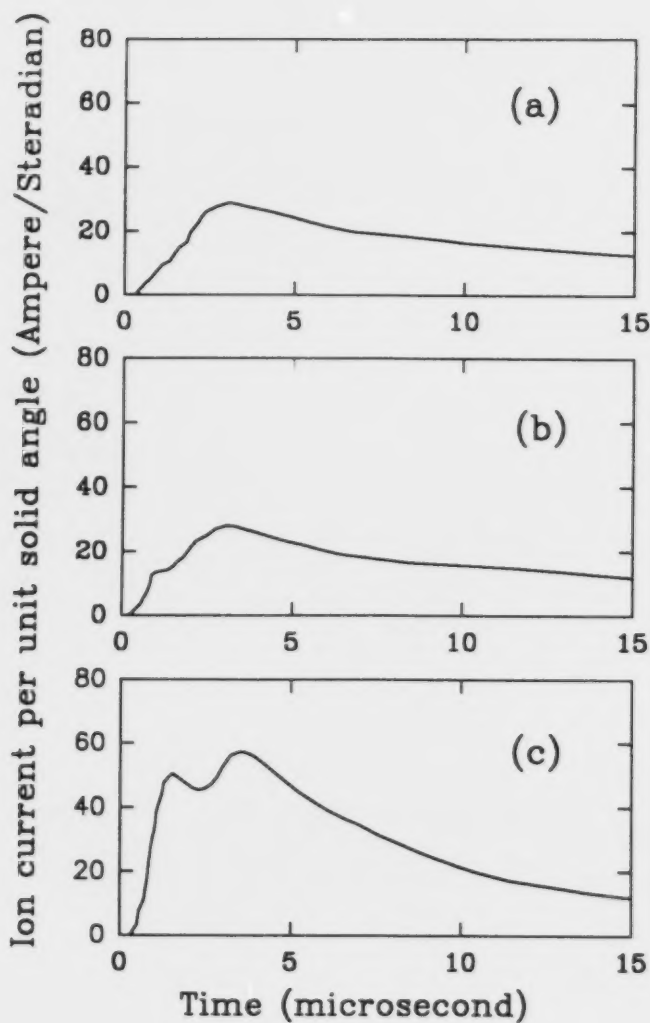


Figure 4.41 Ion current signals of Au plasma observed by a Faraday cup in a direction of 14° with respect to target normal. The (2 ns, $0.268 \mu\text{m}$) laser intensities on target are (a) $1.1 \times 10^{13} \text{ W/cm}^2$ (best-focus), (b) $1.7 \times 10^{12} \text{ W/cm}^2$ (2 mm out-of-focus) and (c) $4.6 \times 10^{11} \text{ W/cm}^2$ (4 mm out-of-focus) respectively.

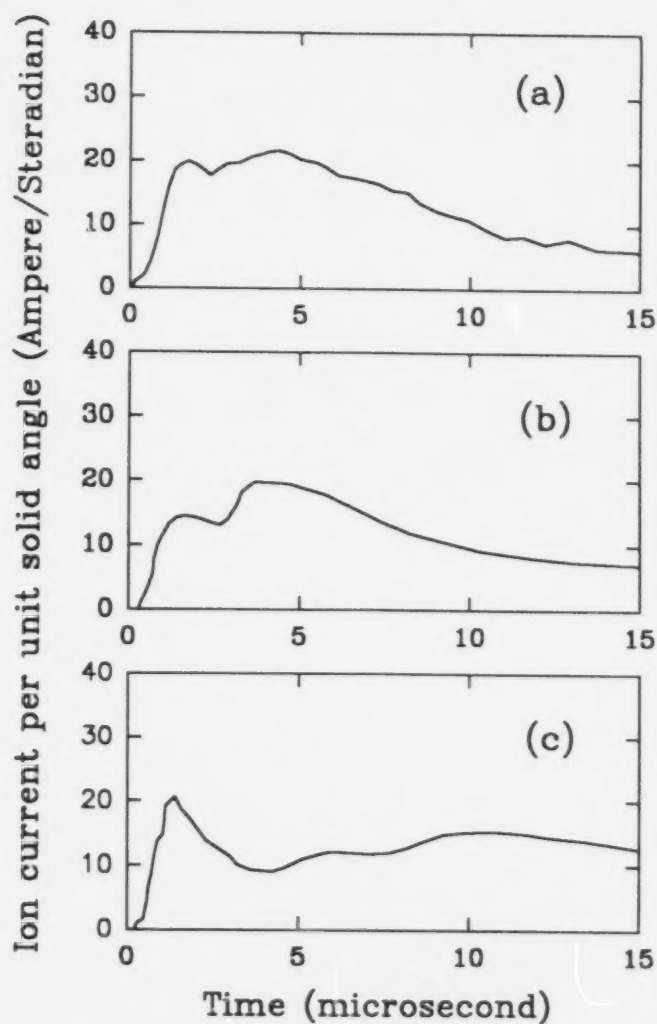


Figure 4.42 Ion current signals of Au plasma observed by a Faraday cup in a direction of 28° with respect to target normal. The (2 ns, $0.268 \mu\text{m}$) laser intensities on target are (a) $1.1 \times 10^{13} \text{ W/cm}^2$ (best-focus), (b) $1.7 \times 10^{12} \text{ W/cm}^2$ (2 mm out-of-focus) and (c) $4.5 \times 10^{11} \text{ W/cm}^2$ (4 mm out-of-focus) respectively.

expected for a nearly planar hydrodynamic expansion (frozen at the end of a laser pulse) resulting in ion blow-off strongly peaked in the direction of target normal.

The above gold ion expansion behavior can be explained⁵⁰ by considering the effect of radiation loss and radiation energy transport using a two-region plasma model consisting of a high temperature electron thermal transport region and a low temperature, high density radiation transport region. The first-peak ions correspond to the hot plasma ablated by electron thermal conduction and the second-peak ions corresponds to the cooler plasma ablated by x-ray radiation. The above picture implies the first-peak ions carry a smaller fraction of the absorbed energy as laser intensity is increased because of the increasing x-ray conversion efficiency as a function of intensity, consistent with experimental observations. In addition, the ions originating from the x-ray radiation transport dominated region should have less angular spread than those from the electron thermal transport dominated region because of their lower temperature again consistent with experimental observations.

4.4 Summary of Asymptotic Behavior of the Expanding Plasma

In general for a fixed direction of observation the magnitude, velocity and average charge state of the ion peak increase with increasing laser intensity for same laser pulse length. Under the same irradiation conditions, as the direction is changed the magnitude, velocity and average charge state of the ion peak decreases with increasing angle.

For short pulse irradiation two peak behavior is observed in aluminum but not carbon. For long pulse irradiation only one single peak is observed in both aluminum and carbon. In general, for the long pulse irradiation because of the lower laser intensity the peak

ion velocity and average charge state of the ion peak is smaller than those for the short pulse irradiation. In addition, for the high intensity irradiation, the carbon ion current and first peak of aluminum ion current are observed to consist of a sharp pulse indicating steady state ablation is taking place. In contrast, for the low intensity irradiation (long pulse and short pulse but large focal spot) the ion currents consist of broad pulses.

The gold ion expansion characteristics are consistent with a two-region plasma model consisting of a high temperature electron thermal transport region and a low temperature, high density radiation transport region.

CHAPTER FIVE

ENERGY BALANCE AND PLASMA CONDITIONS

To understand the physical processes governing the interaction of laser radiation and plasma it is very important to know the fractional absorption of the input laser energy by the targets and the amount of absorbed energy that is converted into radiation and particle (thermal, kinetic and internal) energy.

The most important absorption mechanisms for laser light in a plasma are inverse bremsstrahlung⁵¹ and resonance absorption⁵². The maximum electron density that laser light can reach in a plasma is called the critical density and is given by

$$n_c = 1.1 \times 10^{21} \lambda_L^{-2} \mu\text{m} \quad [\text{cm}^{-3}] \quad (5.1)$$

For KrF laser at $0.248 \mu\text{m}$ this implies the critical density is $1.8 \times 10^{22} \text{ cm}^{-3}$. Resonance absorption can be important if inverse bremsstrahlung absorption does not lead to absorption of light in the underdense region ($n_e < n_c$). For our experiments, at a laser intensity of 10^{13} W/cm^2 , the incident laser light will be attenuated by a factor of two e-foldings by inverse bremsstrahlung absorption at an electron density of $\sim 2 \times 10^{21} \text{ cm}^{-3}$ according to equation (2.66) for an aluminum plasma. At higher laser intensities the laser light can also excite natural modes of oscillation of the plasma by laser-driven plasma instabilities such as stimulated Brillouin scattering (SBS) and stimulated Raman scattering⁵³. These instabilities can lead to scattering of the incident light and thus reduce laser absorption by the plasma. These instabilities require the presence of laser light at densities $n_e \geq 0.1 n_c$ for $I_L \lambda_L^2 \leq 10^{14} \text{ W } \mu\text{m}^2/\text{cm}^2$ for them to grow significantly. The laser light in our experiment only penetrates up to

approximately 1/10 of the critical density implying that these instabilities will not be important in the overall energy balance. The results of the absorption measurements which indicate $> 90\%$ absorption of the laser energy confirm this view. The results of the electron temperature measurements are in good agreement with a steady state ablation model based on inverse bremsstrahlung absorption indicating this process is the dominant absorption mechanism for our experimental conditions. In addition, experiments in this laboratory for similar experimental conditions but at higher intensities ($> 10^{14}$ W/cm²) indicate that only low levels of SBS and negligible levels of SRS are observed⁵⁴.

5.1 Laser Absorption Measurements

The fractional absorption of laser energy by the target can be determined by measuring the incident and scattered light, the difference being the absorbed laser energy. The experimental set up is shown in figure 5.1. The incident laser energy is monitored by a calorimeter. The scattered light measurements consisted of measuring the back reflected light by a calorimeter/beam splitter combination and measuring the diffusely scattered light from the plasma in a 4π geometry by an integrating sphere (see section 3.7). In this experiment a 2 cm diameter aluminum rod is irradiated with 0.268 μ m laser pulses of 2 ns half width. For an average laser intensity of 8×10^{12} W/cm² the average fractional back reflection and diffusively scattered reflection are $0.5 \pm 0.1\%$ and $6.5 \pm 2.0\%$ respectively which resulted in a fractional laser energy absorption of $93 \pm 2\%$.

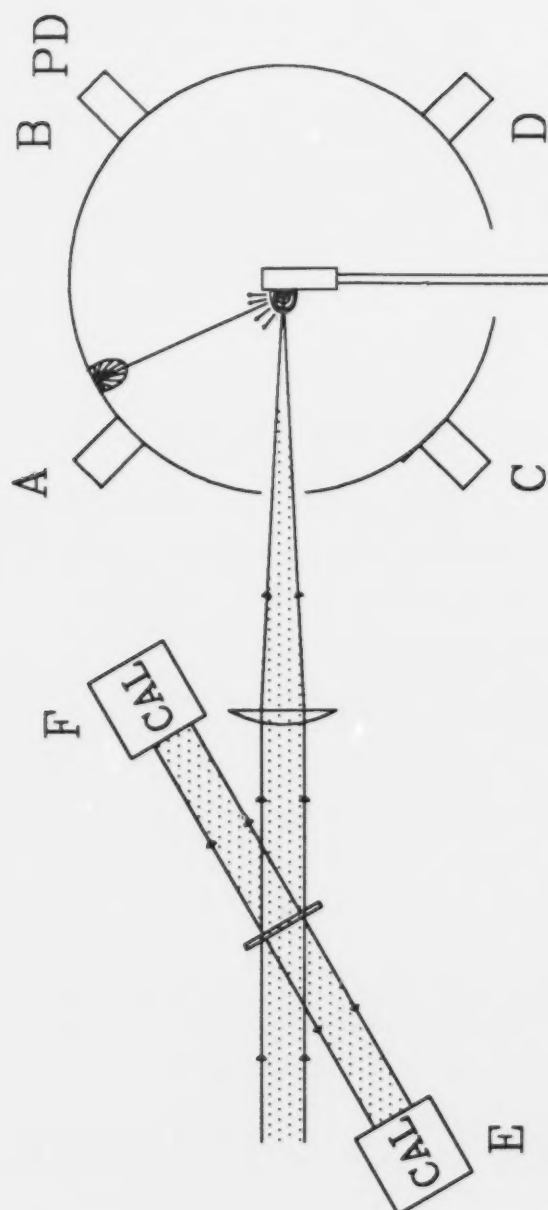


Figure 5.1 Experimental set up using integrating sphere for laser energy absorption measurement: A-D, photodiodes used to measure diffusely scattered light from the plasma; E, calorimeter used to measure the input laser energy; F, calorimeter used to measure back reflected light from the plasma.

5.2 X-Ray Conversion Efficiency

A fast response foil calorimeter (see section 3.9) was used to determine the fraction of absorbed laser energy that was converted into x-rays. The foil calorimeter with a 2 μm aluminum detector element was positioned 6 to 7 cm from the target at an angle of 40° to the target normal to measure the emitted x-ray energy per unit solid angle (see figure 4.1). Since the aluminum absorber foil starts to transmit x-rays of energy > 500 eV, the detector had an effective cutoff energy of 850 eV. In this experiment²⁴ 13 different flat solid or thick foil targets consisting of elements in the range $Z = 6$ to $Z = 82$ were irradiated with a 0.248 μm laser pulse of 2 ns duration (FWHM). The laser intensity was calculated from the 70 % energy-containing area of the focal spot which corresponded to an average equivalent diameter of 100 μm . For an average laser intensity of 5×10^{12} W/cm² the soft x-ray energies as a function of atomic number observed at an angle of 40° relative to the target normal were measured. The integrated x-ray conversion efficiency can be obtained if the angular distribution of the x-ray emission is known. The angular x-ray emission should follow a $(\cos \theta)^\delta$ law with δ lying between 0 and 1. The conversion efficiency of KrF laser energy into soft x-rays, in the photon energy range of $\sim 10 - 850$ eV, is presented in figure 5.2 as a function of atomic number for various target materials. The results are plotted as full circles with error bars assuming a $\cos \theta$ ($\delta = 1$) emitting distribution and as crosses, assuming an isotropic distribution, into 2π steradians ($\delta = 0$). Assuming a $\cos \theta$ distribution, the conversion efficiency was determined to be 6 ± 2 % for carbon ($Z = 6$), 14 ± 4 % for aluminum ($Z = 13$) and 31 ± 8 % for gold ($Z = 79$) targets.

Using arrays of soft x-ray diodes with similar experimental parameters, measurements by R. Popil⁵⁵ have

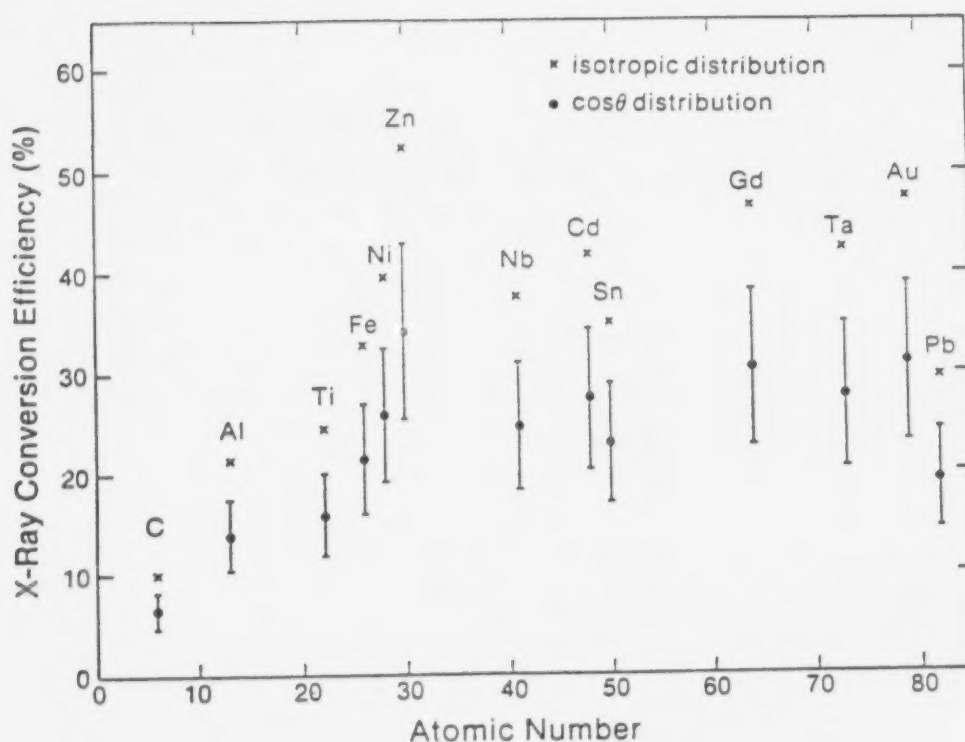


Figure 5.2 Soft X-ray conversion efficiency ($h\nu < 850$ eV) as a function of Z for $I = 5 \times 10^{12}$ W/cm² for cases of a $\cos\theta$ and an isotropic distribution. The error bars on the isotropic distribution points have been omitted for clarity.

indicated that for aluminum and gold the x-ray emission distributions obey a $(\cos \theta)^{0.76}$ and $(\cos \theta)^{0.33}$ laws respectively. These results are expected since gold is a more optically thick x-ray emitter than aluminum and thus the angular emission distribution should be more isotropic. Incorporating these angular corrections lead to an increase of x-ray conversion efficiencies of aluminum and gold by factors of 1.07 and 1.26 respectively. The resulting 10 - 850 eV x-ray conversion efficiencies for aluminum and gold are $15 \pm 4 \%$ and $39 \pm 10 \%$ respectively.

The total x-ray conversion efficiencies can be calculated from an estimation of the fraction of x-ray energy transmitted through the $2 \mu\text{m}$ aluminum detector element. Under the present irradiation conditions, a temperature of approximately 200 eV is expected for the gold plasma. About a quarter of the radiation from a 200 eV black body will be transmitted through a $2 \mu\text{m}$ aluminum foil (see figure 5.3). A black-body spectrum represents the upper limit for the radiation from optically thick lines which will dominate the real emission spectrum of the elements. Thus total conversion efficiencies for aluminum and gold of approximately 20 % and 50 % can be estimated. These results are in good agreement with the 15% and 54%^{55,56} x-ray conversion efficiencies obtained for aluminum and gold respectively using an array of soft x-ray diodes and x-ray streak camera under similar experimental conditions.

Our Al conversion efficiency agrees reasonably well with other data for UV radiation at $\sim 10^{14} \text{ W cm}^{-2}$ which gave 20 %⁵⁷ with a measured angular distribution and 6 - 12 %^{58,59} assuming a $\cos \theta$ angular distribution law. Our Au conversion efficiency also agree reasonably well with the value of 36 % ($I_L = 2 \times 10^{14} \text{ W cm}^{-2}$, $\tau_L = 0.45 \text{ ns}$, $D_{f.s.} = 50 \mu\text{m}$) assuming a $\cos \theta$ angular distribution law reported in reference 59. But our data is somewhat lower than other reported values of 72 %⁶⁰ ($I_L = 1 \times 10^{14} \text{ W cm}^{-2}$, $\tau_L = 1 \text{ ns}$,

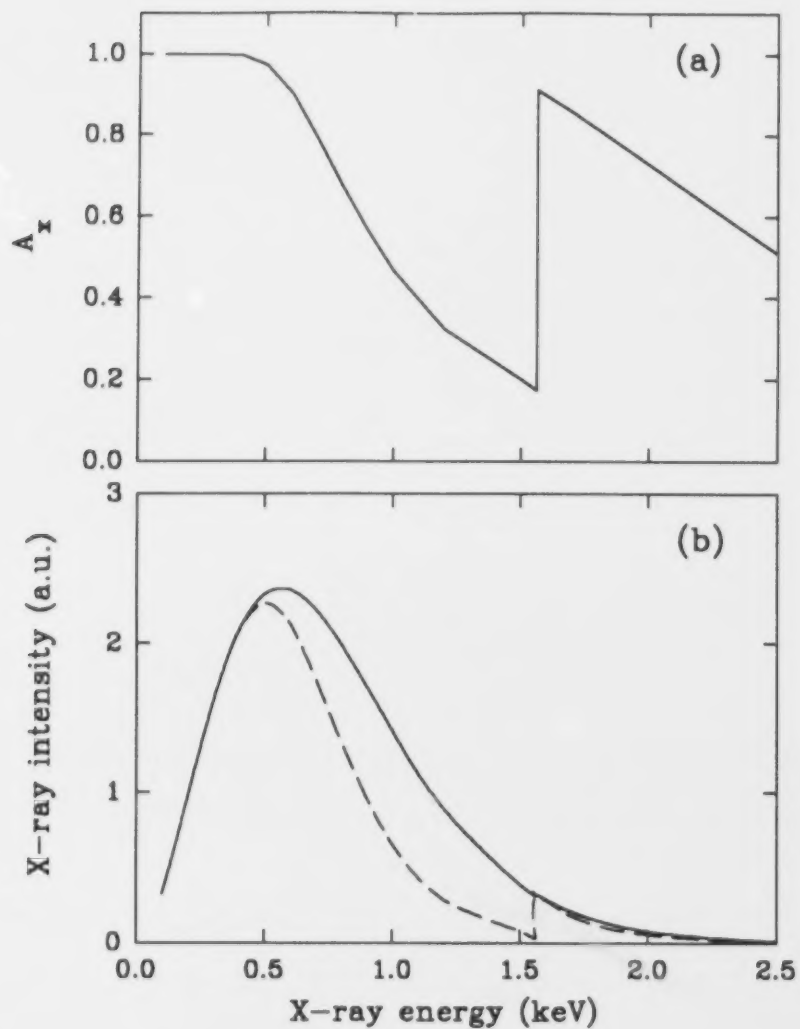


Figure 5.3 (a) Fraction of X-ray energy absorbed, A_x , as a function of X-ray photon energy for a 2 μm Al foil; (b) X-ray intensity spectrum radiated by a 200 eV blackbody (solid line) and X-ray intensity spectrum absorbed by a 2 μm Al foil (dash line).

$D_{f.s.} = 1400 \text{ } \mu\text{m}$) and 80%⁵⁷ ($I_L = 2.4 \times 10^{14} \text{ W cm}^{-2}$, $\tau_L = 0.4 \text{ ns}$, $D_{f.s.} = 150 \text{ } \mu\text{m}$). This may be due to the smaller focal spot and longer pulse duration used in our experiment.

5.3 Thermal and Kinetic Energy of the Expanding Plasma

The absorbed laser energy produces a plasma expanding outward. Initially most of the absorbed laser energy is converted to thermal energy, ionization energy and other internal energy of the plasma. For 1 J of absorbed laser radiation at best focus with $I_L = 1 \times 10^{13} \text{ W/cm}^2$ the expected mass ablation rate is $1.1 \times 10^5 \text{ g cm}^{-2} \text{ s}^{-1}$ (see equation 2.59). For a 2 ns laser pulse with a focal spot diameter of $74 \text{ } \mu\text{m}$ the mass ablated is approximately 10 ng. This implies approximately 20 mJ of the energy is invested in ionizing the plasma since the $\approx 500 \text{ eV}$ plasma produced at this irradiance should be He-like ($Z = 11$). Thus most of the absorbed laser energy becomes thermal energy of the plasma. As the plasma expands, part of this energy will be emitted as photons and a large fraction of the remaining energy will become kinetic energy of the ions with some remaining as internal and thermal energy of the plasma.

An array of plasma calorimeters (see figure 4.1) is used to measure the plasma energy distribution at different observation angles of the expanding plasma. The plasma calorimeter, in principle, measures the total energy of the plasma, as given by the following equation

$$E_{cal} = E_{xray} + E_{K.E.} + E_{ioniz} + E_{SL} \quad (5.2)$$

where E_{xray} , $E_{K.E.}$, E_{ioniz} and E_{SL} are x-ray energy, kinetic energy of ions, ionization energy of plasma and energy of scattered light respectively.

An array of charge collectors is used to determine the angular distribution of the ion kinetic energy of the expanding plasma as discussed in chapter 4. The charge

collectors measure the angular ion current density $I(\theta, t)$ as a function of observation angle, θ , which is

$$I(\theta, t) = Z(\theta, t) e \frac{d}{dt} [n_i(\theta, t)] \quad (5.3)$$

where $n_i(\theta, t)$ and $Z(\theta, t)$ are the number density and charge state of the ions as a function of angle. The angular distribution of ion kinetic energy can be obtained from

$$E_{k.E.}(\theta) = \int \frac{A m_p v^2(\theta, t) I(\theta, t)}{e Z(\theta, t)} dt \quad (5.3)$$

Thus ion collector measurements and a knowledge of the charge state distribution of the plasma (see chapter 4) can yield the kinetic energy distribution.

Figure 5.4 and 5.5 summarize the energy per unit solid angle measured by plasma calorimeters and ion charge collectors for carbon and aluminum plasma, respectively, produced by 0.268 μm and 0.248 μm laser pulses with half width of 2 ns at various laser intensities. A plume model^{61,62} developed originally for a rocket exhausting into a vacuum is used to fit the data. According to the model the energy per unit solid angle E_θ can be expressed as a function of the angle θ from the target normal by

$$E_\theta = E_0 \cos^\alpha(\beta \theta) \quad (5.4)$$

with

$$\alpha = \frac{2}{\gamma - 1} \quad (5.5)$$

where E_0 is the energy per unit solid angle at $\theta = 0^\circ$, γ the heat capacity ratio and β is a parameter which depends on the Mach number M of the plasma at the cross section where the transition of planar flow to spherical flow takes

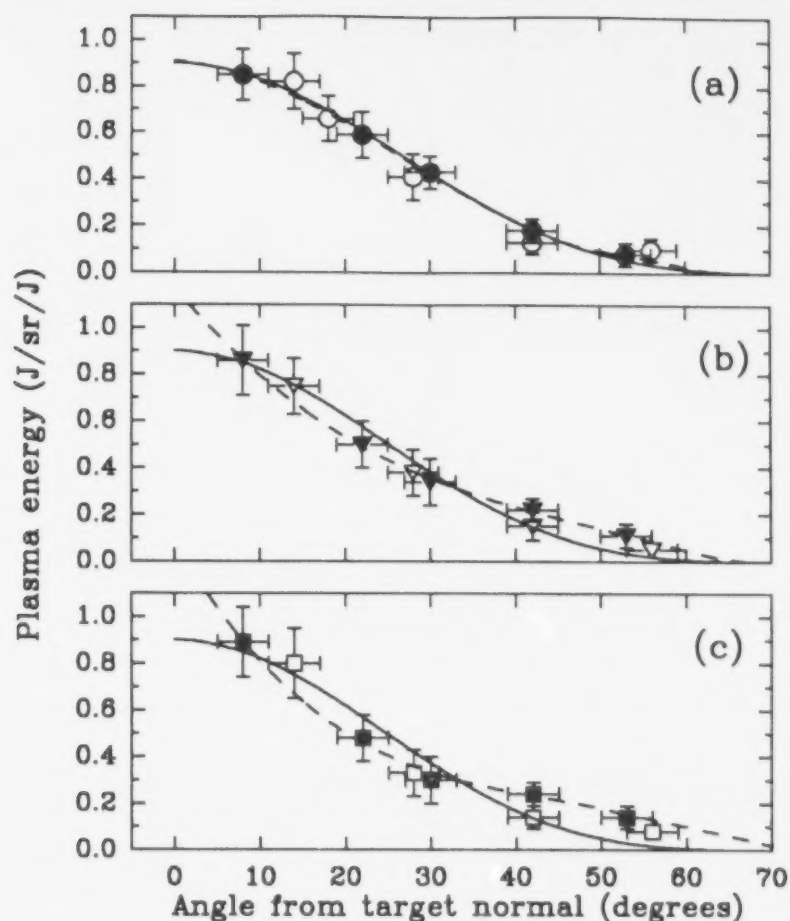


Figure 5.4 Angular distributions of the plasma energy for C plasma produced by $0.248 \mu\text{m}$, 2 ns laser pulses at different irradiation conditions: (a) best-focus, $I_L = 8 \times 10^{12} \text{ W/cm}^2$, (b) 2 mm off-focus, $I_L = 3 \times 10^{12} \text{ W/cm}^2$, (c) 4 mm off-focus, $I_L = 1 \times 10^{12} \text{ W/cm}^2$. The solid and hollow symbols represent plasma calorimeter data and ion charge collector data respectively. The solid curve and dashed curves are best fits to the calorimeter data using the plume model and a 3rd order polynomial fit respectively.

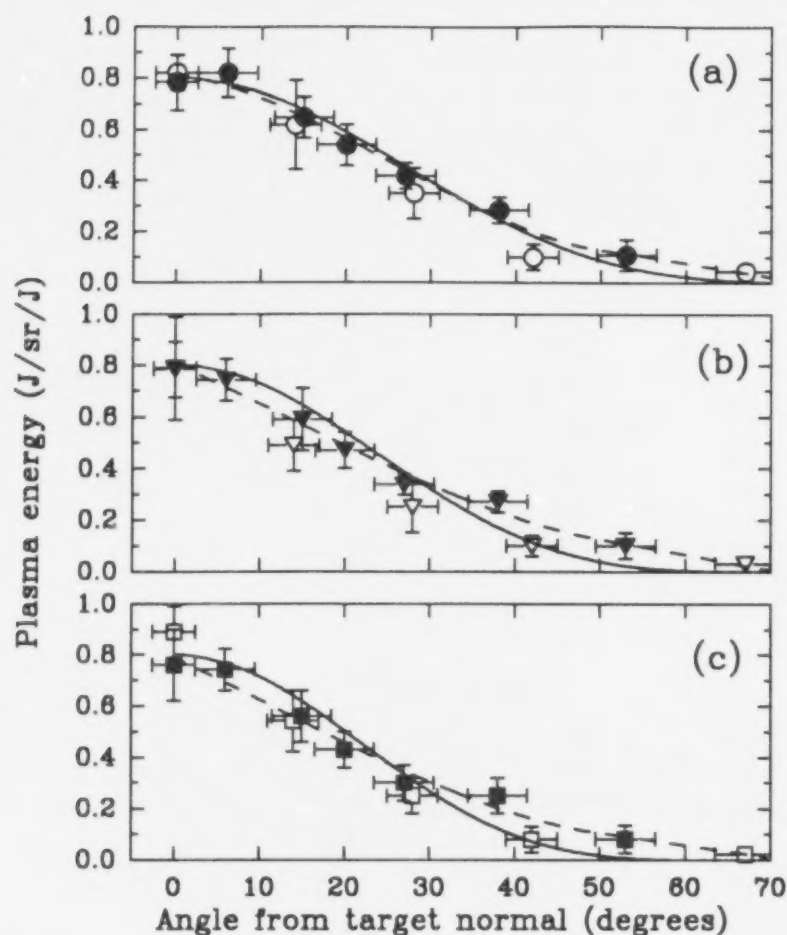


Figure 5.5 Angular distributions of the plasma energy for Al plasma produced by 2 ns, 0.248 μm and 0.268 μm laser pulses at different irradiation conditions: (a) best-focus, $I_L = 8 \times 10^{12} \text{ W/cm}^2$, (b) 2 mm off-focus, $I_L = 3 \times 10^{12} \text{ W/cm}^2$, (c) 4 mm off-focus, $I_L = 1 \times 10^{12} \text{ W/cm}^2$. The solid and hollow symbols represent plasma calorimeter data and ion charge collector data respectively. The solid curve and dashed curves are best fits to the calorimeter data using the plume model and a 3rd order polynomial fit respectively.

place. For $\gamma = 1.5$ the angular dependence of energy per steradian is expected to vary as $\cos^4(\beta \theta)$. For best focus (f0) irradiation the agreement between the model (see solid curves) and the data is excellent. For off-focus (f1 and f2) irradiation, the model (solid curve) fits the data very well for observation angles less than 30 degrees but the data falls off less rapidly than the model predicts for angle larger than 30 degrees. Table 5.1 summarizes the parameter β and the resulting Mach number M obtained by a least squares fit to the plasma calorimeter data.

Table 5.1 Parameter β obtained from fitting the plasma calorimeter data for different laser irradiation conditions and the resulting Mach number M .

	f0		f1		f2	
	C	Al	C	Al	C	Al
β	1.13	1.10	1.21	1.24	1.22	1.35
M	2.3	2.2	2.6	2.7	2.6	3.0

For best focus irradiation, $M \approx 2$ fits both the carbon and aluminum data. Figure 5.6 gives the predicted Mach number on laser axis of an aluminum plasma produced by an $0.248 \mu\text{m}$ laser pulse of 1.5 ns FWHM as a function of the ratio of distance, d , from the target surface to the diameter of focal spot, $d_{f.s.}$, as calculated by the 2D hydro code Castor. The peak laser intensity on target is $2 \times 10^{13} \text{ W/cm}^2$ and the focal spot diameter is $68 \mu\text{m}$ in this case. According to the Castor result, $M \approx 2$ at the region $d/d_{f.s.} = 1 - 2$ where the transition from planar to spherical flow takes place; this is in agreement with the plume model. The total energy can be obtained by integrating $E(\theta)$. For best focus irradiation condition the

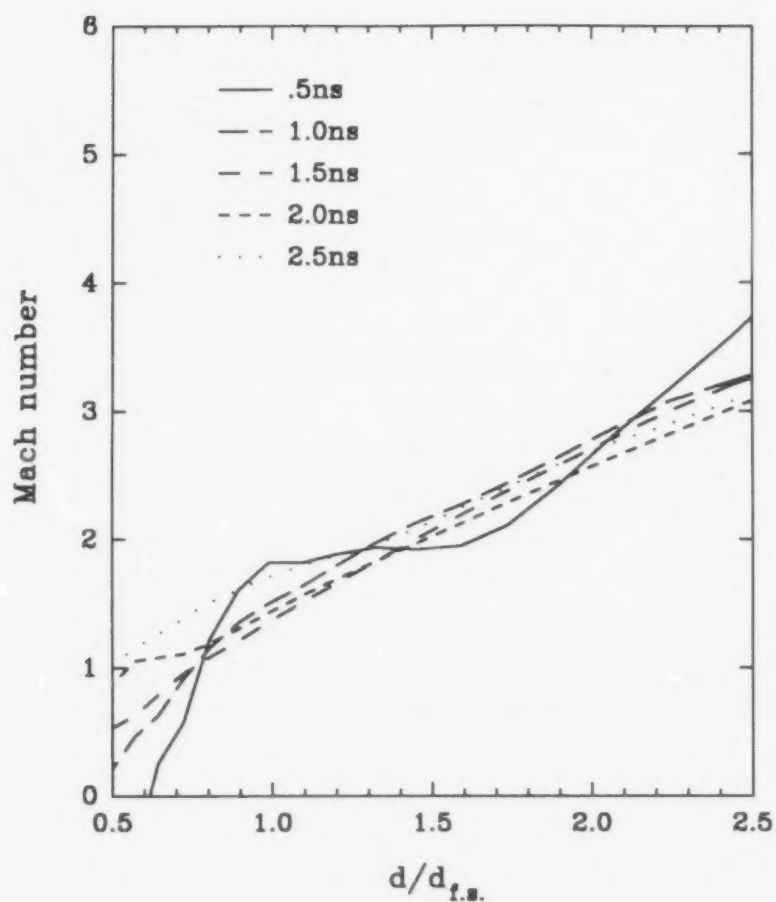


Figure 5.6 Mach number (M) on laser axis of an aluminum plasma produced by an $0.248 \mu\text{m}$ laser pulse of 1.5 ns FWHM as a function of the ratio of distance, d , from the target surface to the diameter of focal spot, $d_{f.s.}$, as calculated by the 2D hydro code Castor. The peak laser intensity on target is $2 \times 10^{13} \text{ W/cm}^2$ and the focal spot diameter is $68 \mu\text{m}$.

area under the solid curves in figure 5.4 (a) and 5.5 (a) give 0.89 ± 0.15 J/J and 0.84 ± 0.15 J/J for C and Al plasma respectively. For off-focus irradiation conditions the solid curves are used from $\theta = 0^\circ$ to 30° and the dashed curves are used for $\theta > 30^\circ$. Results are summarized in column D of Table 5.2 in the following section.

The kinetic energy per unit solid angle as measured by ion charge collectors are plotted as hollow symbols in figure 5.4 and figure 5.5. For carbon and aluminum plasma, the total energy of the plasma as measured by plasma calorimeters and kinetic energy of ions as measured by ion charge collectors, in general, agree with each other within 20 %. This indicates most of the input laser energy becomes kinetic energy of the ions at distances large from the target surface. Similar results are observed for carbon and aluminum plasma produced by 20 ns 0.248 μm and 0.694 μm laser pulses.

In case of gold plasma, the calorimeter energy accounted for 85 ± 20 % of the incident laser energy. A comparison of the calorimetric data with ion charge collector data can be used to distinguish between the ion and x-ray components of the measured energy. The ion-current signals are analyzed by using the measured charge dependence as a function of velocity for the case of a tungsten plasma⁶³ as an approximation for the gold plasma. The ratio of kinetic energy of ions to the calorimeter energy measured in the direction of $\theta = 14^\circ$ as a function of laser energy is shown in figure 5.7. It is seen that this ratio decreases with increasing laser intensity and is smallest for the high-intensity best focus irradiation. In comparison, for C and Al plasma this ratio is close to unity irrespective of laser energy and focal spot size. The decreasing ratio of kinetic energy of ions to the calorimeter energy clearly implies an increasing component of radiation energy emitted from the plasma, in agreement with the x-ray emission energy results.

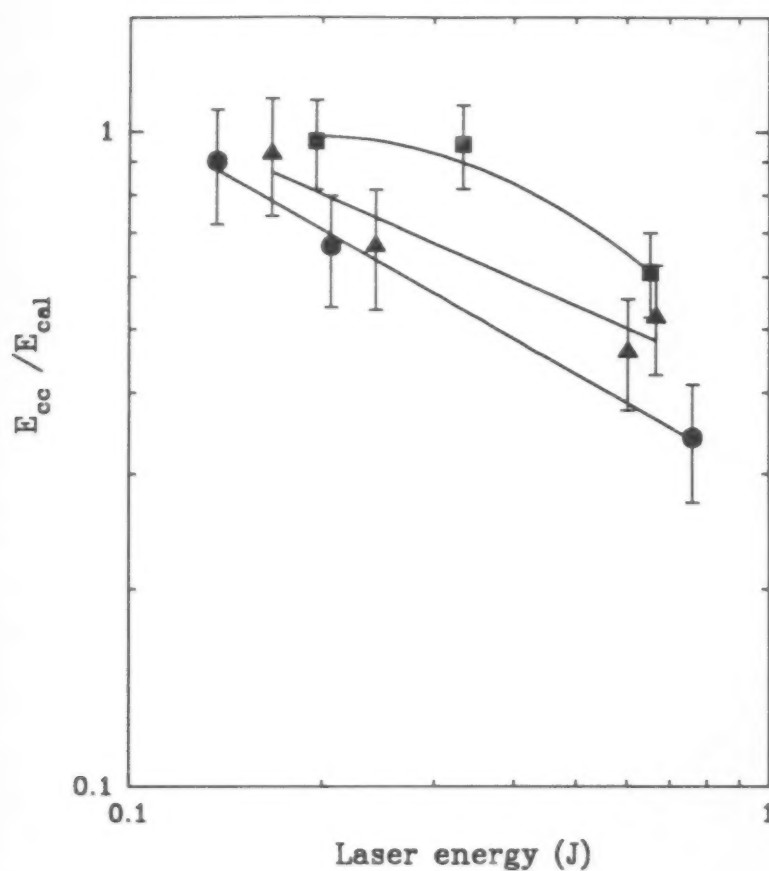


Figure 5.7 Dependence of the ratio of ion kinetic energy (calculated from ion charge collectors) to the plasma energy (measured by calorimeters) ($\theta = 14^\circ$) on laser energy and focal spot for gold plasma. The circles, triangles and squares represent best-focus data, 2 mm off-focus data and 4 mm off-focus data respectively.

5.4 Summary of energy balance

The measured energy balance of carbon, aluminum and gold plasma produced by different laser irradiation conditions are summarized in Tables 5.2 - 5.4. Column A gives the type of plasma; column B gives the focal condition of the laser beam; column C gives the laser intensities on target; column D gives the total energy of the plasma as measured by the plasma calorimeters; column E gives an estimate of laser energy invested in ionizing the plasma by assuming a coronal equilibrium condition; column F gives the x-ray energy from the plasma as measured by the x-ray calorimeter; column G gives the kinetic energy of the ions as calculated from the ion charge collector signals.

Table 5.2 Energy Balance for plasma produced by 2 ns KrF laser pulse.

A	B	C	D	E	F	G
Plasma Condition			Energy Distribution (Fraction of absorbed energy)			
	focus cond.	Intensity (W/cm ²)	Total	Ioniz	X-ray	Kinetic
C	f0	8×10^{12}	.89±.15	.04	.05±.02	.85±.25
	f1	3×10^{12}	.94±.15	.07		.78±.25
	f2	1×10^{12}	.98±.15	.12		.83±.25
Al	f0	6×10^{12}	.84±.20	.02	.20±.07	.79±.25
	f1	1.5×10^{12}	.86±.20	.04		.71±.25
	f2	3×10^{11}	.78±.20	.07		.67±.25
Au	f0	5×10^{12}	.88±.25	.11	.50±.12	.35±.15
	f1	1×10^{12}	.86±.25			.69±.25
	f2	3×10^{11}	.83±.25			.78±.25

Table 5.3 Energy Balance for plasma produced by 20 ns KrF laser pulse.

A	B	C	D	E	F	G
Plasma Condition			Energy Distribution (Fraction of absorbed energy)			
	focus cond.	Intensity (W/cm ²)	Total	Ioniz	X-ray	Kinetic
C	f0	2x10 ¹¹	.91±.20	.06		.83±.25
Al	f0	2x10 ¹¹	.93±.20	.10		.79±.25

Table 5.4 Energy Balance for plasma produced by 25ns Ruby laser pulse.

A	B	C	D	E	F	G
Plasma Condition			Energy Distribution (Fraction of absorbed energy)			
	focus cond.	Intensity (W/cm ²)	Total	Ioniz	X-ray	Kinetic
C	f0	2x10 ¹¹	.91±.20	.06		.88±.25
Al	f0	2x10 ¹¹	.87±.20	.10		.84±.25

5.5 Plasma Conditions

The two parameters that describe the thermodynamic

state of the plasma are temperature and density. The detailed spatial and temporal evolution of these parameters will be determined primarily from the 2D hydrodynamic code Castor. The validity of temperature results are checked by experimental measurements.

5.5.a Plasma Density Profile

The laser produced plasma will expand away from the target surface in a two dimensional plasma flow. However, the initial expansion of the plasma will be approximately one dimensional planar as long as its forward expansion remains less than the lateral radius of the plasma. Since the plasma expands at approximately the ion acoustic velocity C_s the duration for which the 1D planar expansion should be valid is $\approx d_{fs}/C_s$ where d_{fs} is the laser focal spot diameter. For an aluminum plasma produced by a 70 μm diameter laser beam at an irradiance of $1 \times 10^{13} \text{ W/cm}^2$, the maximum plasma temperature should be equal to $\approx 500 \text{ eV}$. This results in a maximum ion acoustic velocity of $\approx 1.4 \times 10^7 \text{ cm/s}$. This implies the plasma flow should be approximately planar for the first half nanosecond of the laser pulse. As the size of the plasma grows, the initial plasma region will approach a point source compared to the expanding plasma and eventually the plasma flow can be approximated by a one dimensional spherical flow model.

The time evolution of electron number density contours and plasma flow vectors from a 2D hydrodynamic simulation Castor are given in figures 5.8 to 5.11. The aluminum plasma is created by focusing a 0.248 μm triangular laser pulse, with a peak intensity of $I_0 = 2 \times 10^{13} \text{ W/cm}^2$, a rise time of 600 ps and FWHM of 1.5 ns. The time origin at $t = 0$ corresponds to the start of the triangular pulse. The plasma flow is approximately planar at early times and becomes approximately spherical at late times.

Interferometric measurements have been made in a

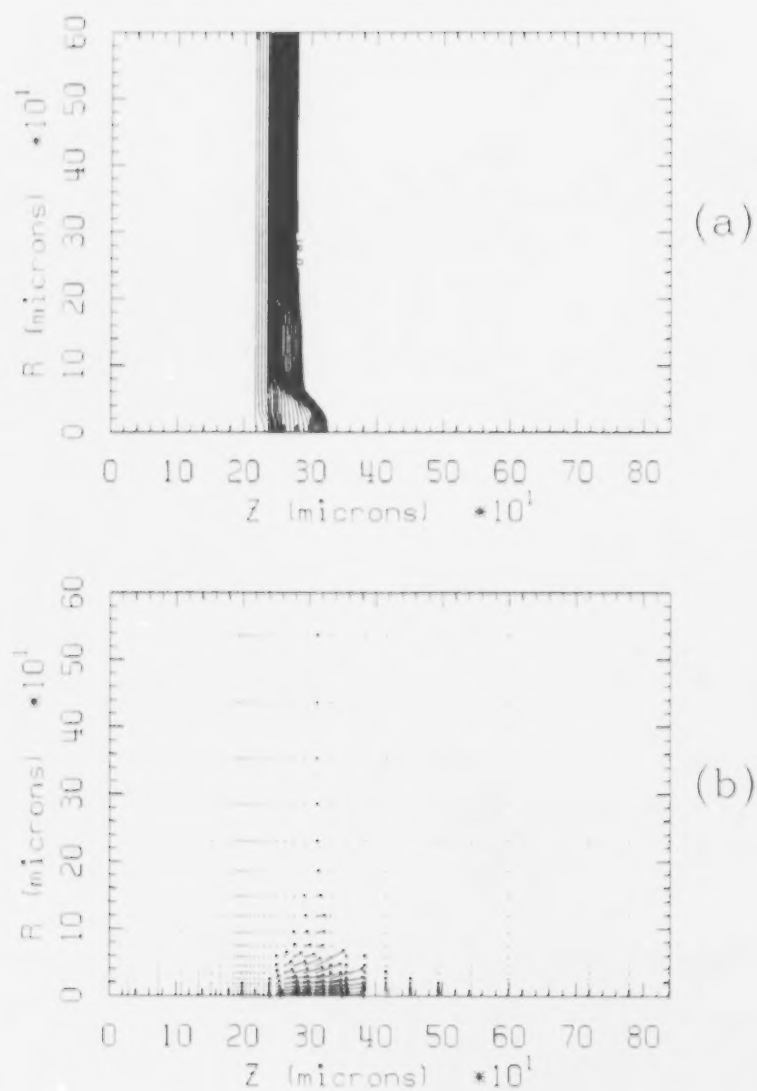


Figure 5.8 2D Castor simulation results for an aluminum plasma $I_0 = 2 \times 10^{13}$ W/cm², $\tau_L = 1.5$ ns at $t = 0.25$ ns: (a) electron density contour (log scale), (b) vector flow plot.

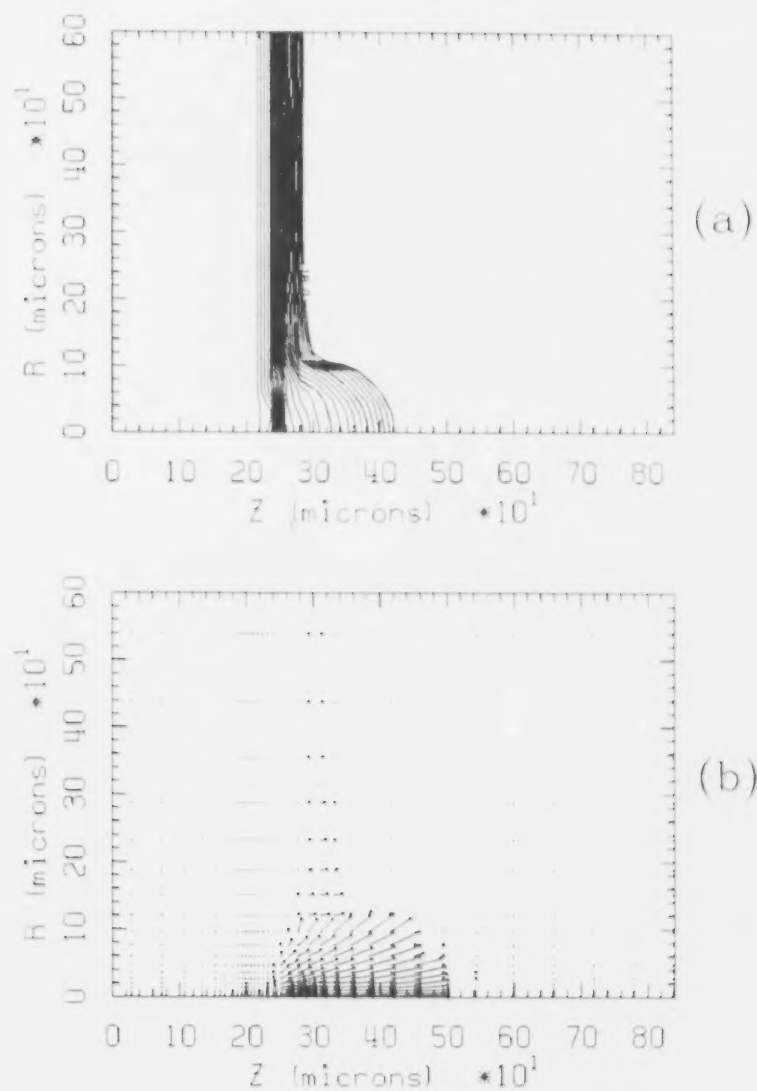


Figure 5.9 2D Castor simulation results for an aluminum plasma $I_0 = 2 \times 10^{13}$ W/cm², $\tau_L = 1.5$ ns at $t = 0.50$ ns: (a) electron density contour (log scale), (b) vector flow plot.

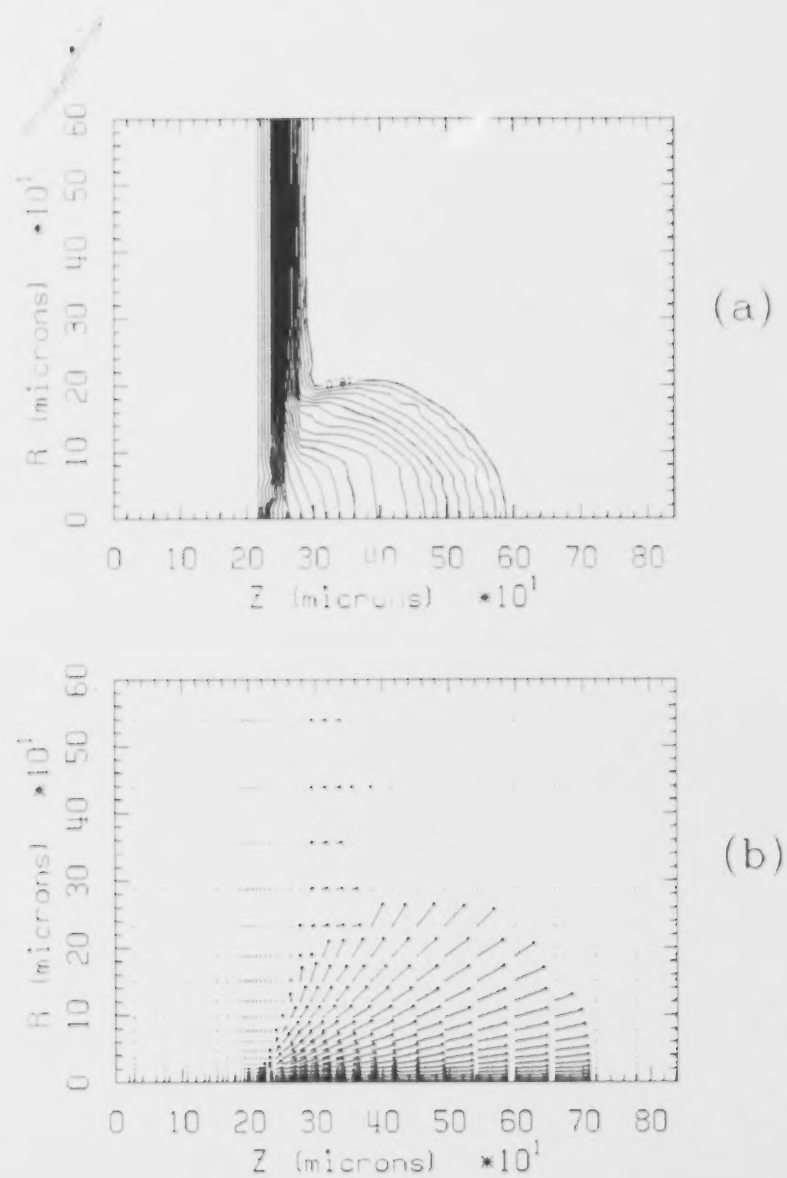


Figure 5.10 2D Castor simulation results for an aluminum plasma $I_0 = 2 \times 10^{13}$ W/cm², $\tau_L = 1.5$ ns at $t = 1.0$ ns: (a) electron density contour (log scale), (b) vector flow plot.

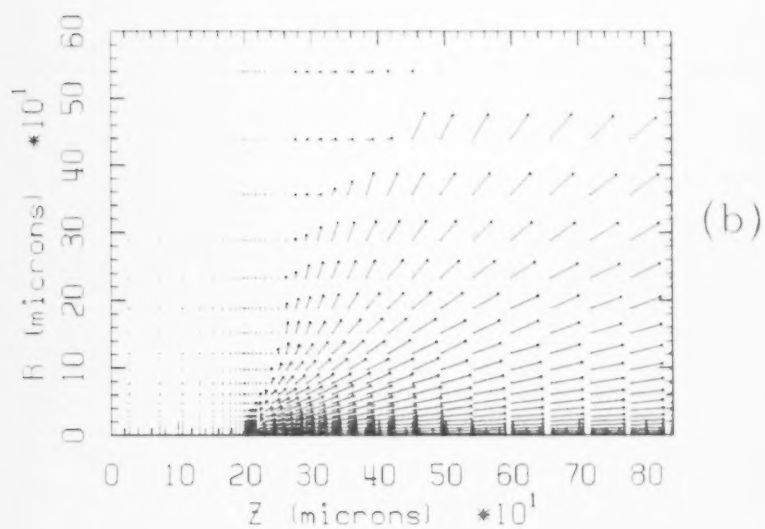
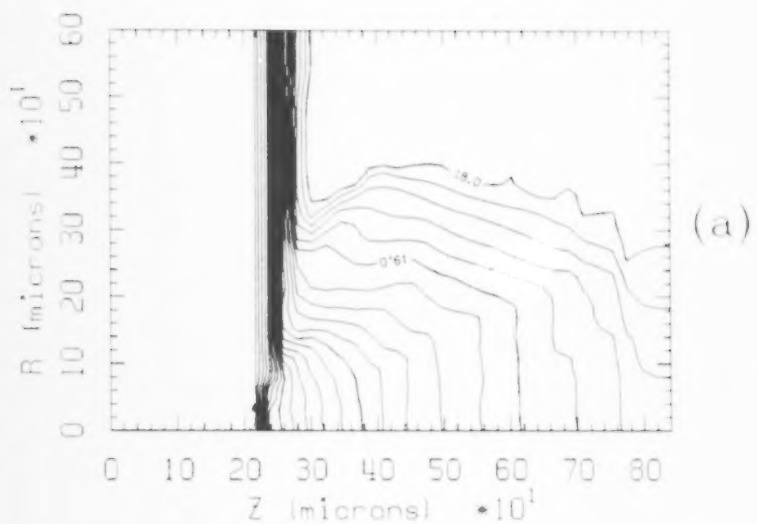


Figure 5.11 2D Castor simulation results for an aluminum plasma $I_0 = 2 \times 10^{13}$ W/cm², $\tau_L = 1.5$ ns at $t = 2.0$ ns: (a) electron density contour (log scale) (b) vector flow plot.

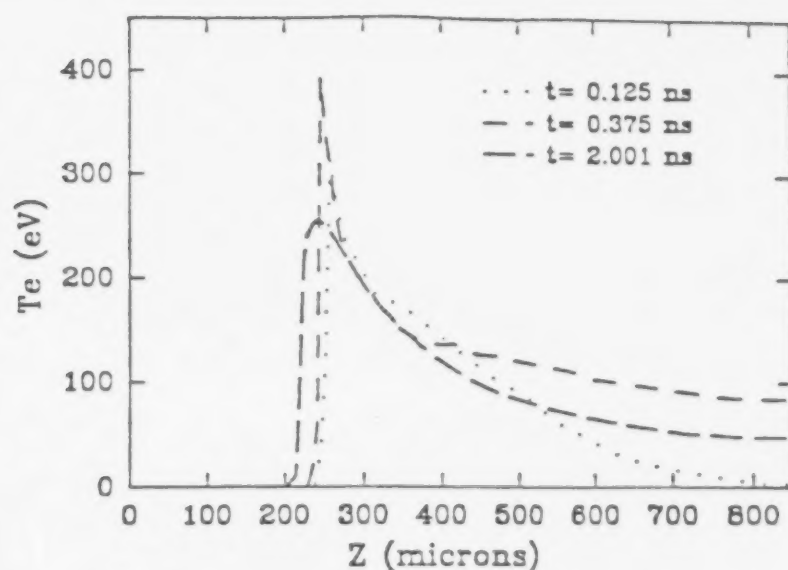
separate investigation¹⁴ using a nitrogen laser to study aluminum plasma created by focusing a 1.5 ns KrF laser beam to a 35 μm focal spot diameter. Those results show that for the first nanosecond of the laser pulse the initial plasma on the laser axis is characterized by an approximately planar expansion. A smooth transition from planar to stationary spherical flow takes place after the first nanosecond of the laser pulse. The results are, in general, in good agreement with simulation results from the 2D hydrodynamic code Castor except for some fine details (possibly due to the lack of magnetic fields, radiation transport and time dependent ionization physics in the simulation).

5.5.b Plasma Temperature Profile

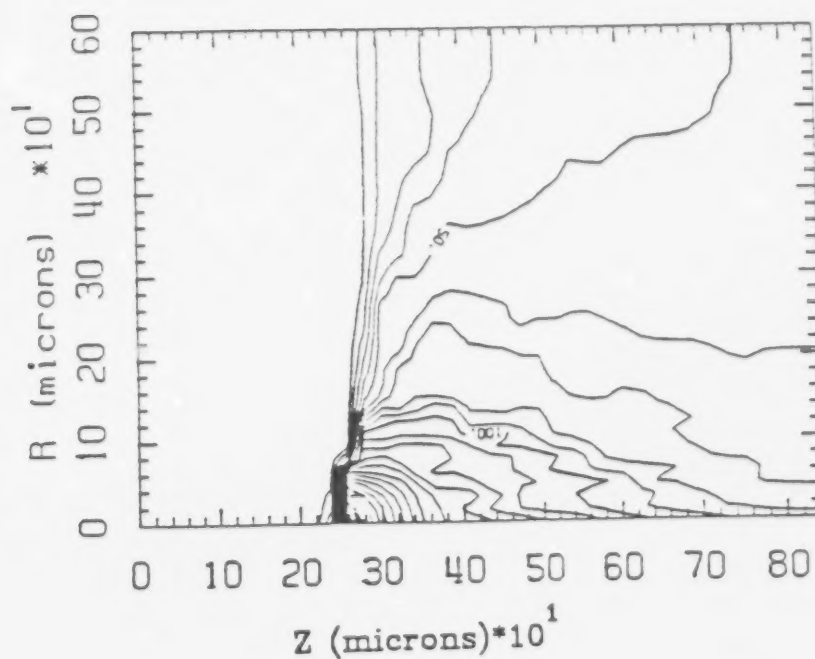
The electron temperature distribution profiles from a 2D hydrodynamic simulation Castor are given in figure 5.12 for an aluminum plasma. The plasma behavior is calculated assuming a 0.248 μm triangular laser pulse with a peak intensity of $I_0 = 2 \times 10^{13} \text{ W/cm}^2$, a rise time of 600 ps and FWHM of 1.5 ns. The temperature profile very quickly takes the shape shown for time 0.375 ns which corresponds to the simulation's peak temperature. The maximum temperature at the laser focus is 400 eV.

5.5.b.1 Determination of Plasma Temperature from X-Ray Emission

The temperature of a laser plasma can be obtained by using a metal foil absorber scintillator technique (see section 3.10). Since this technique is based on the measurement of high energy x-rays ($>1 \text{ keV}$) and these x-rays are most efficiently emitted during the period when the plasma is at its peak temperature, this measurement will be weighed heavily toward the temperature maximum.



(a)



(b)

Figure 5.12 2D Castor simulation results for an Al plasma $I_0 = 2 \times 10^{13}$ W/cm², $\tau_L = 1.5$ ns at $t = 2.0$ ns: (a) electron temperature profile for $t = 0.125$ ns (dot), $t = 0.375$ ns (medium dash), $t = 2.0$ ns (long dash); (b) electron temperature contour plot at $t = 0.375$ ns.

5.5.b.2 Determination of Plasma Temperature from Ion Velocities

The average temperature of the plasma can be determined from the measurement of the peak ion velocity. The assumption of steady state ablation (see section 2.5.c) implies a near delta function ion distribution is produced. Since most of the thermal energy of the initial plasma is converted into kinetic energy of the expanding plasma at a distance where the ion distribution is measured, a measurement of the kinetic energy of the expanding plasma should reveal the temperature of the initial plasma where the ablation took place.

The internal energy for a multiply ionized plasma is given by

$$\rho \epsilon = \frac{3}{2} p + n_i I \quad (5.6)$$

$$I = \sum_{i=1}^Z E_{\infty, i} \quad (5.7)$$

since for $T_e = T_i$

$$p = (1 + Z) n_i k T_e \quad (5.8)$$

it follows

$$\rho \epsilon = p \left(\frac{3}{2} + \frac{I}{(1+Z)k T} \right) \quad (5.9)$$

For an ideal gas the quantity, γ , defined as the ratio of heat capacity at constant volume to heat capacity at constant pressure is equal to 5/3. In analogy to this constant heat capacity ratio one can define an effective heat capacity ratio γ' for a plasma by⁶³

$$\rho \epsilon = \frac{p}{\gamma' - 1} \quad (5.10)$$

It follows from equation (5.9) that

$$\gamma' = 1 + \frac{1}{\frac{3}{2} + \frac{I}{(1+Z)kT}} \quad (5.11)$$

For an adiabatic spherical expansion where the plasma flow is stationary there exists a boundary separating the upstream subsonic flow and downstream supersonic flow. The plasma velocity is equal to the acoustic velocity at this boundary. The total energy of the plasma at the acoustic cross section is given by

$$E_0 = \rho_0 \epsilon_0 + p_0 + \frac{1}{2} \rho_0 C_{sa}^2 \quad (5.12)$$

where the adiabatic acoustic speed is given by

$$C_{sa} = \sqrt{\gamma' \frac{p_0}{\rho_0}} \quad (5.13)$$

The total energy of the plasma downstream that is at a distance far away from the target surface where most of its thermal energy is converted to kinetic energy is given by

$$E_\infty = \rho_\infty \epsilon_\infty + p_\infty + \frac{1}{2} \rho_\infty v_\infty^2 \quad (5.14)$$

Since the flow is stationary it follows from the continuity equation and energy conservation equation that

$$\frac{E_0}{\rho_0} = \frac{E_\infty}{\rho_\infty} \quad (5.15)$$

If one ignores the pressure downstream which is small compared with the kinetic energy density, one has

$$\epsilon_0 + \frac{p_0}{\rho_0} + \frac{1}{2} C_{sa}^2 = \epsilon_\infty + \frac{1}{2} v_\infty^2 \quad (5.16)$$

One can consider two extreme cases: (a) recombination is very large so most of the ionization energy is returned to the plasma; (b) recombination is small and most of the ionization energy still remains as internal energy in the plasma. For case (a) $\epsilon_\infty \approx 0$, it follows from equations (5.14), (5.11) and (5.8) that

$$v_\infty^2 = \frac{\gamma' + 1}{\gamma' - 1} C_{sa}^2 \quad (5.17)$$

For case (b) $\epsilon_\infty = n_i I_0$, the internal energy downstream is equal to the ionization energy upstream, it follows from equations (5.16), (5.6), (5.10) and (5.13) that

$$v_\infty^2 = \frac{\gamma' + 5}{\gamma'} C_{sa}^2 \quad (5.18)$$

Equations (5.17) and (5.18) together with equations (5.8) and (5.13) express the upstream temperature as a function of the downstream velocity for case (a) and (b) respectively as

$$T_e = \frac{A}{Z+1} \frac{m_p}{e} \frac{\gamma' - 1}{\gamma' + 1} \frac{v_\infty^2}{\gamma'} \quad (5.19)$$

and

$$T_e = \frac{A}{Z+1} \frac{m_p}{e} \frac{v_\infty^2}{\gamma' + 5} \quad (5.20)$$

The temperature of the plasma can be determined if the charge state of the plasma is known. For simplicity, We will assume the charge distribution of the plasma can be described by a coronal model. The exact charge distribution is not very important since γ' is only a slowly varying function, lying between 1.2 to 1.6 for a wide range of Z and T values as shown in figure 5.13. In Figure 5.14 the temperature as a function of ion velocity is shown for case (a) (see equation 5.19) and case (b) (see equation 5.20). In the real situation where a fraction of the initial ionization energy is returned to the plasma, the true temperature of the plasma should lie between the two limits given by cases (a) and (b).

5.5.b.3 Summary of Temperature Measurements

The measurements of plasma temperature for aluminum, carbon and gold plasmas produced by a ≈ 2 ns KrF laser pulse under various focal conditions are summarized in figures 5.15 to 5.17 respectively. The graphs include results from x-ray emission; for aluminum and carbon plasma they also include results from velocities of peak ion current based on equations (5.19) and (5.20) (dotted lines). Temperature scaling using the steady state ablation model based on equations (2.58) and (2.62) are shown as dashed lines. Temperature scaling as a function of laser intensity from the x-ray emission data are found by a least square fit to the emission data and are plotted as solid lines in figures 5.15 to 5.17. The equations for these lines are given by

$$\text{CARBON: } T_e = 439 I_{13}^{0.56} \quad [\text{eV}] \quad (5.21)$$

$$\text{ALUMINUM: } T_e = 515 I_{13}^{0.58} \quad [\text{eV}] \quad (5.22)$$

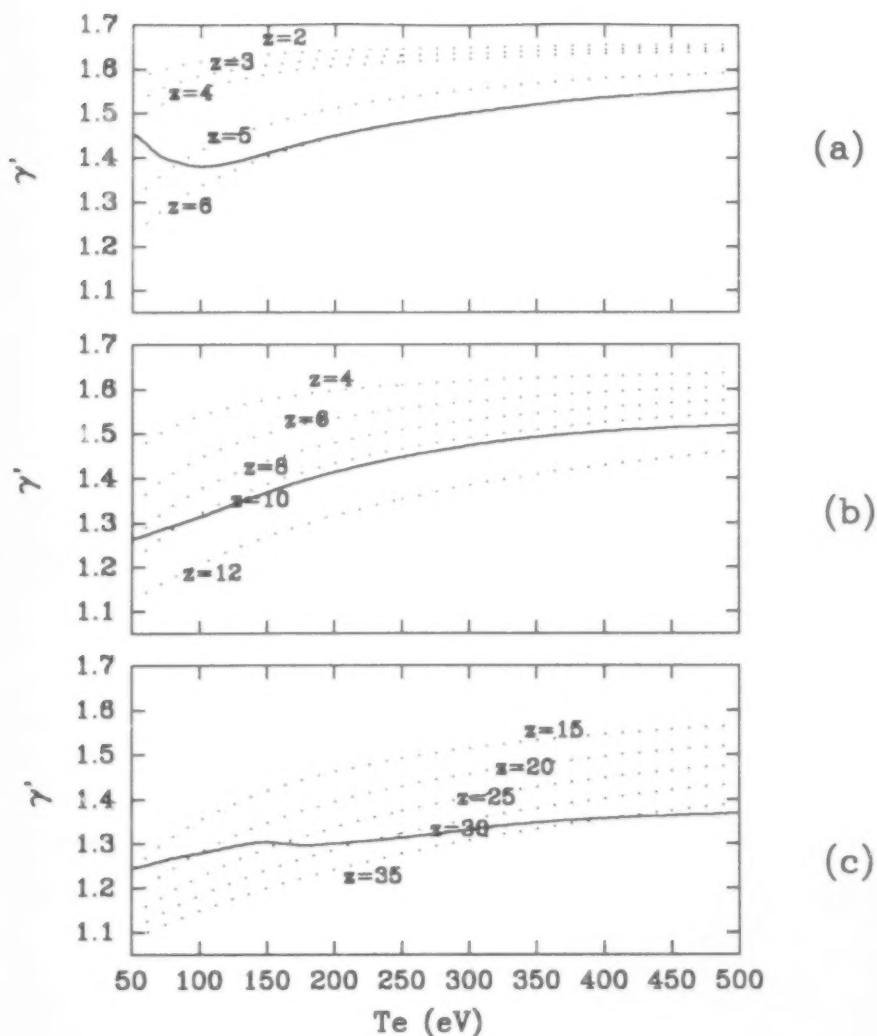


Figure 5.13 Effective heat capacity ratio (γ') as a function of temperature for (a) carbon plasma (b) aluminum plasma (c) gold plasma. The different dotted curves represent γ' for different fixed value of Z . The solid curve is the calculated value where Z is obtained from the coronal equilibrium model.

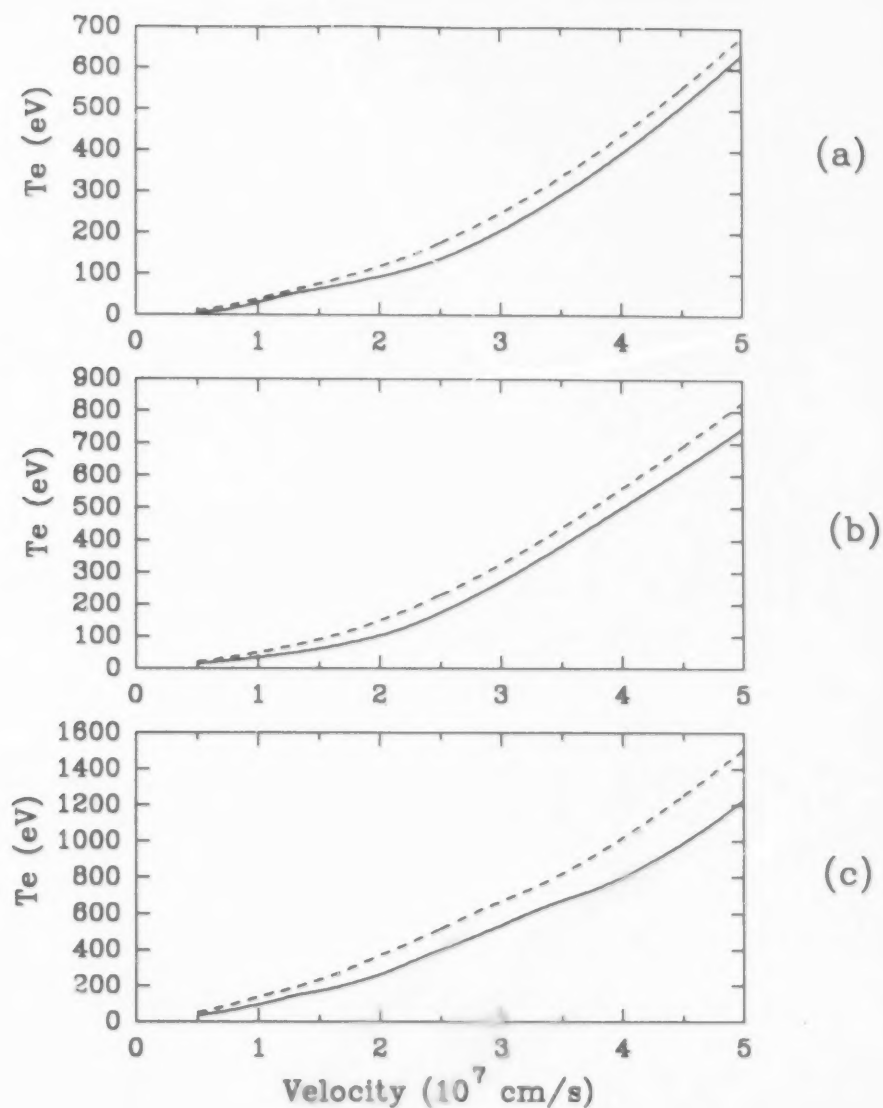


Figure 5.14 Temperature as a function of ion velocity for (a) carbon plasma (b) aluminum plasma (c) gold plasma. The solid and dashed curves represent cases when maximum and no recombination is assumed respectively.

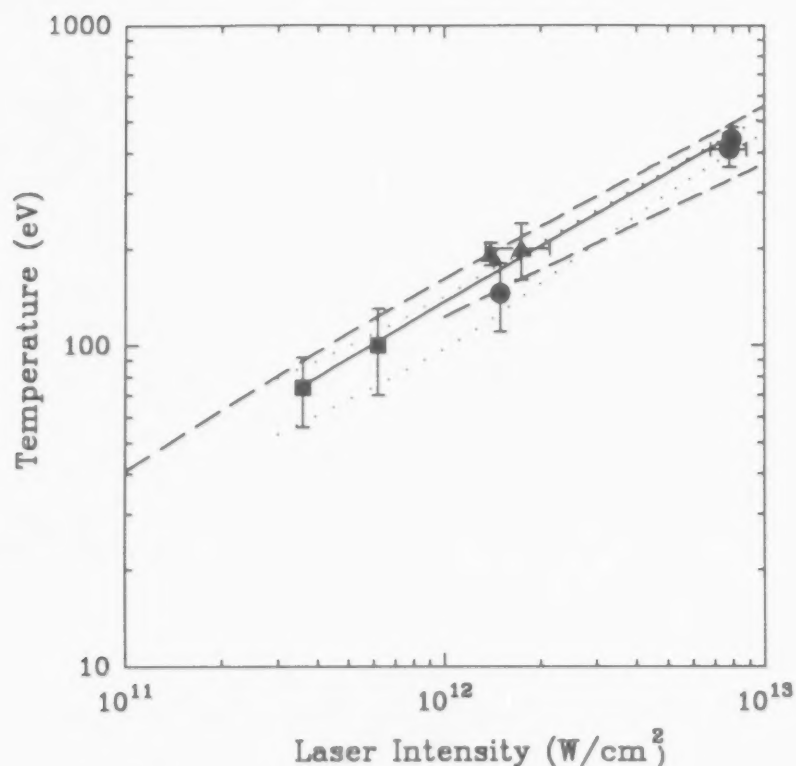


Figure 5.15 Electron temperature of an aluminum plasma as a function of laser intensity (2 ns, 0.248 μm & 0.268 μm): results from x-ray emission with the circles, triangles and squares representing data from the best focus (f0), 2 mm off focus (f1) and 4 mm off focus (f2) irradiation conditions; results from ion velocities based on equations (5.18) and (5.17) (upper and lower dotted lines) and temperature scaling from steady state ablation model based on equations (2.58) and (2.62) (upper and lower dashed lines).

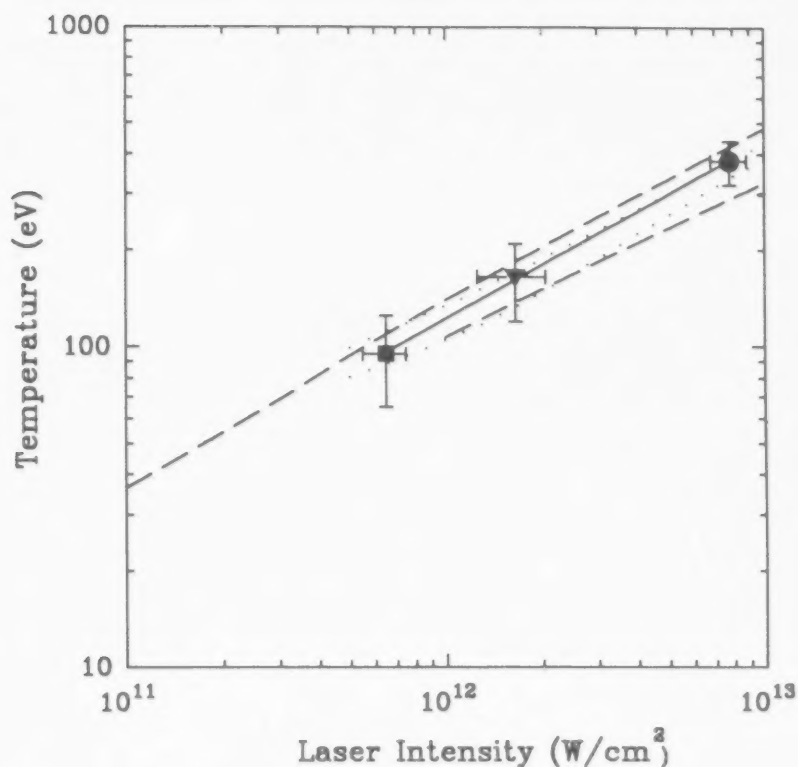


Figure 5.16 Electron temperature of a carbon plasma as a function of laser intensity (2 ns, 0.248 μm): results from x-ray emission with the circles, triangles and squares representing data from the best focus (f0), 2mm off focus (f1) and 4mm off focus (f2) irradiation conditions; results from ion velocities based on equations (5.18) and (5.17) (upper and lower dotted lines) and temperature scaling from steady state ablation model based on equations (2.58) and (2.62) (upper and lower dashed lines).

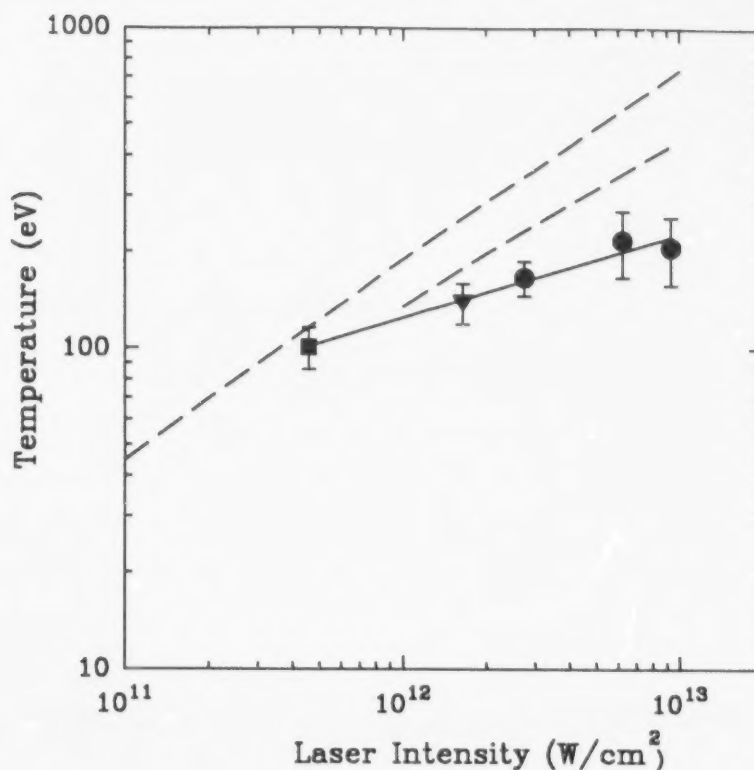


Figure 5.17 Electron temperature of a gold plasma as a function of laser intensity (2 ns, 0.268 μm): results from x-ray emission with the circles, triangles and squares representing data from the best focus (f_0), 2 mm off focus (f_1) and 4 mm off focus (f_2) irradiation conditions; temperature scaling from steady state ablation model based on equations (2.58) and (2.62) (upper and lower dashed lines).

$$\text{GOLD: } T_e = 231 I_{13}^{0.26} \quad [\text{eV}] \quad (5.23)$$

For carbon and aluminum plasma a scaling of $T_e \sim I^{0.56-0.58}$ is observed which is higher than that predicted by the steady state ablation model where $T_e \sim I^{0.44}$. However, the measured temperatures are lower than those predicted by the model at the smaller intensity. This is consistent with the fact that a greater fraction of the absorbed energy is invested in ionizing the plasma in the f2 focus condition than the f0 focus condition (see table 5.2). The larger fraction of energy invested in ionizing the plasma implies less energy is available for heating the plasma. The steady state model can be modified by calculating the effective energy available for the plasma by subtracting the energy invested in ionizing the plasma from the laser energy absorbed by the plasma

$$I_{\text{eff}} = \frac{E_a - E_{\text{ionz}}}{\pi r_{fs}^2 \tau_L} \quad (5.21)$$

where E_a , E_{ionz} , r_{fs} , and τ_L are the laser energy absorbed by the plasma, energy invested in ionizing the plasma, the focal spot radius and the pulse length of the laser. The charge state distribution required to calculate E_{ionz} is calculated by assuming a coronal equilibrium model. For laser intensities in the range 10^{11} W/cm^2 to 10^{13} W/cm^2 the modified temperature scaling laws are, from equations (5.21) and (2.63)

$$\text{CARBON: } T_e = 445 I_{13}^{0.57} \quad [\text{eV}] \quad (5.24)$$

$$\text{ALUMINUM: } T_e = 541 I_{13}^{0.57} \quad [\text{eV}] \quad (5.25)$$

Tables 5.5 and 5.6 summarize the plasma temperature

measurements from x-ray emission (column C) and peak ion velocity assuming maximum recombination heating and no recombination heating (column D) and; results from the steady state ablation model (column E) and the modified steady state model (column F) for carbon and aluminum plasmas produced by various laser conditions.

Table 5.5 Temperature measurements and results from models for 2 ns KrF laser pulse (C plasma)

A	B	C	D	E	F	G
PLASMA CONDITIONS		ELECTRON TEMPERATURE (eV)				
			Peak Ion Velocities		Steady State Ablation Model	
	I_L (W/cm ²)	X-ray	Maximum Recomb.	No Recomb.	Normal	Modified
f0	1×10^{13}	439	388	434	480	445
f1	2×10^{12}	179	148	188	204	177
f2	5×10^{11}	83	81	100	94	81

Table 5.6 Temperature measurements and results from models
for 2 ns KrF laser pulse (Al plasma)

A	B	C	D	E	F	G
PLASMA CONDITIONS		ELECTRON TEMPERATURE (eV)				
	I_L (W/cm ²)	X-ray	Peak Ion Velocities		Steady State Ablation Model	
			Maximum Recomb.	No Recomb.	Normal	Modified
f0	1×10^{13}	515	461	518	558	541
f1	2×10^{12}	203	155	209	234	216
f2	5×10^{11}	91	68	100	109	98

Table 5.7 Temperature measurements and results from models
for 20 ns KrF laser pulse (C plasma)

A	B	C	D	E	F	G
PLASMA CONDITIONS		ELECTRON TEMPERATURE (eV)				
	I_L (W/cm ²)	X-ray	Peak Ion Velocities		Steady State Ablation Model	
			Maximum Recomb.	No Recomb.	Normal	Modified
f0	1.6×10^{11}	45	44	53	42	41

Table 5.8 Temperature measurements and results from models
for 20 ns KrF laser pulse (Al plasma)

A	B	C	D	E	F	G
PLASMA CONDITIONS		ELECTRON TEMPERATURE (eV)				
	$I_L (W/cm^2)$	X-ray	Peak Ion Velocities		Steady State Ablation Model	
			Maximum Recomb.	No Recomb.	Normal	Modified
f0	1.6×10^{11}	40	35	51	47	45

Table 5.9 Temperature measurements and results from models
for 25 ns ruby laser pulse (C plasma)

A	B	C	D	E	F	G
PLASMA CONDITIONS		ELECTRON TEMPERATURE (eV)				
	$I_L (W/cm^2)$	X-ray	Peak Ion Velocities		Steady State Ablation Model	
			Maximum Recomb.	No Recomb.	Normal	Modified
f0	1.0×10^{11}	50	42	51	58	48

Table 5.10 Temperature measurements and results from models
for 25 ns ruby laser pulse (Al plasma)

A	B	C	D	E	F	G
PLASMA CONDITIONS		ELECTRON TEMPERATURE (eV)				
	I_L (W/cm ²)	X-ray	Peak Ion Velocities		Steady State Ablation Model	
			Maximum Recomb.	No Recomb.	Normal	Modified
f0	1.0×10^{11}	45	44	48	67	54

The plasma temperature measurements from x-ray emission and from ion velocities are in good agreement for all cases. In general, at short pulse high irradiance, the plasma temperature measurements from x-ray emission agree better with those from ion velocities when no recombination is assumed. At long pulse low irradiance, the x-ray emission temperature falls in between the ion velocity temperatures. This is somewhat expected since recombination is more important for plasma created by lower laser irradiance where the temperature is lower.

Temperature measurements are in reasonable agreement with the modified temperature scaling law for all cases. For short pulse high irradiation of aluminum, the remaining discrepancy of about 6 % in temperatures from x-ray measurements and the modified model can be accounted for by the additional radiation energy loss which is not included in the modified model. In addition, for $I_L \approx 8 \times 10^{12}$ W/cm² the measured temperature from x-ray emission is equal to 450 ev which is in good agreement with the peak temperature

of 400 eV from the Castor simulation for similar irradiation condition.

For gold plasma, the temperature scaling is observed to be $T_e \propto I^{0.26}$, which is much slower than that predicted by the steady state ablation model. At high laser intensity ($\sim 10^{13} \text{ W/cm}^2$) the measured temperature is considerably lower than the predicted values of 400 - 700 eV from the model. In addition, for this laser intensity the measured temperatures for gold plasma is considerably smaller than those of carbon and aluminum plasma, contrary to what is expected from the steady state ablation model. These results can be explained by the fact that over half the absorbed energy is radiated away in x-rays and that the plasma hydrodynamics is dominated by x-ray radiation energy transport.

5.6 Summary of Energy Balance and Plasma Conditions

Absorption of laser energy is found to be $> 90 \%$. The sum of measured x-ray energy, kinetic energy and thermal energy of ions approximately equals the absorbed laser energy. Reasonable agreement between the experimental angular distribution of plasma energy and that predicted by a plume model indicates applicability of the plume model.

The plasma temperature and its intensity scaling deduced from x-ray continuum emission and ion diagnostics agree reasonably well with a steady state ablation model based on inverse bremsstrahlung when the effect of energy invested in ionizing the plasma is included.

The plasma hydrodynamics for gold plasma is dominated by x-ray radiation energy transport.

CHAPTER SIX

ABLATION PARAMETERS AND ENERGY TRANSPORT

Mass ablation rate measurements are sensitive indicators of heat transport processes and the ablation pressure governs the hydrodynamic response of the target.

Spatially and temporally integrated ablation parameters can be determined from measurements of the asymptotic ion flow using charge collectors and time integrated plasma energy using plasma calorimeters.

Ablation parameters⁴⁹ are deduced from measurements of ion emissions produced by laser pulses of typically 2 ns duration at 0.27 μm focused onto planar aluminum targets to produce irradiances between 10^{11} - 10^{13} W/cm^2 . Three different procedures are used in the evaluation of experimental data. The results are compared with results from a steady state analytical model and 1D and 2D numerical calculations. Finally, the results are compared with those from other laboratories.

6.1 Evaluation Procedures for Experimental Data

Spatially and temporally integrated values of mass ablation rate ($\dot{m}_a$) and ablation pressure (P_a) can be obtained from the values of total ablated mass (M) and the momentum ($p_\perp$) carried off perpendicular to the target surface via the following equations

$$\dot{m}_a = \frac{M}{A_a \tau_a} \quad (6.1)$$

$$P_a = \frac{P_\perp}{A_a \tau_a} \quad (6.2)$$

where A_a and τ_a are the ablation area and time respectively. Since the mass ablation rate is expected to scale as $I^{5/9}$ for a spherically steady state ablation model (see section 2.5.c), the ablation time τ_a is taken as the FWHM of the laser pulse scaled by a power of 5/9. A typical laser pulse used in the experiments can be approximated by a 1.5 ns triangular pulse with a rise time of 600 ps. For such pulses the ablation time is $\tau_a = 2.4$ ns (see figure 6.1). The ablation area A_a is taken as the area containing 90 % of the total laser energy. M and $p_{\perp}$ can be calculated from the ion emission data measured by charge collectors and plasma calorimeters. For each expansion angle, the ion charge collector data gives the ion current as a function of time and the plasma calorimeter data gives the plasma energy. After expanding to a large distance most of the plasma energy has been converted into kinetic energy hence

$$E = \frac{1}{2} M \langle v^2 \rangle \quad (6.3)$$

where E is the total plasma energy and $\langle v^2 \rangle$ the mean square velocity of the expanding plasma. Thus

$$M = \frac{2 E}{\langle v^2 \rangle} \quad (6.4)$$

$$p = M \langle v \rangle \quad (6.5)$$

The calculations of ablated mass and momentum can be carried out using three different procedures.

6.1.a Procedure 1

The plasma energy per unit solid angle $E_{cal}(\theta)$ from plasma calorimeter measurements and the peak angular ion velocity $v_p(\theta)$ from the charge collector measurements are

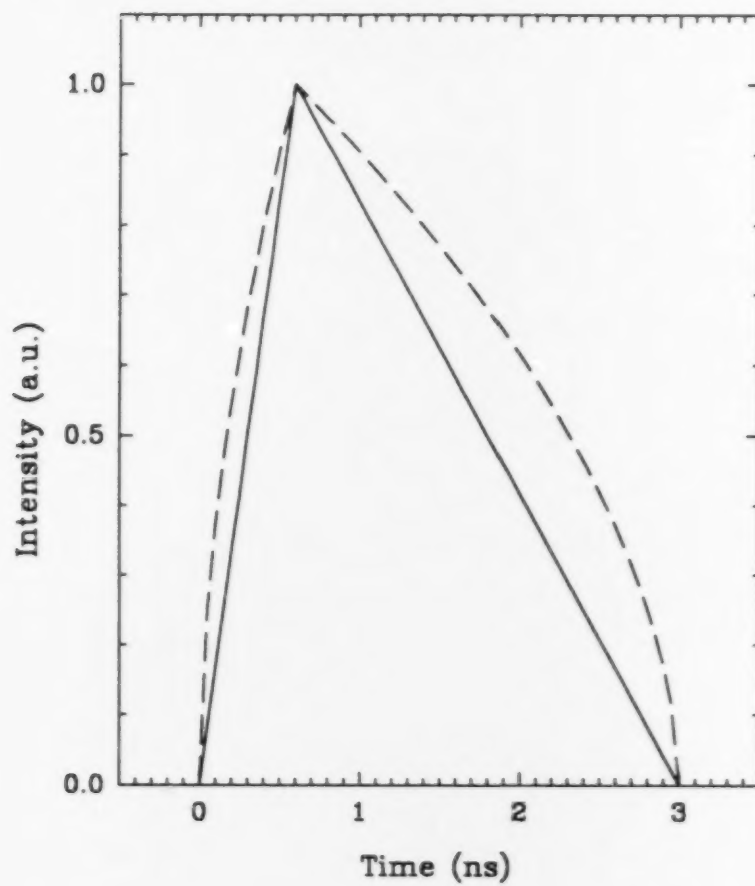


Figure 6.1 Schematic diagram of a typical laser pulse (solid line) and the pulse to the 5/9 power (dashed line).

used to deduce the quantities by integrating over the solid angle

$$M = \int \frac{2 E_{cal}(\theta)}{v_p^2(\theta)} d\Omega \quad (6.6)$$

$$p_{\perp} = \int \frac{2 E_{cal}(\theta) \cos(\theta)}{v_p(\theta)} d\Omega \quad (6.7)$$

The $\cos(\theta)$ factor in equation (6.7) is due to the fact that only the component of the plasma momentum in the direction of the target normal contributes to the net momentum transfer to the target.

6.1.b Procedure 2

Since the charge collectors measure ion current $i(t)$ in a given solid angle, the total mass of the ablated ions can be determined if the ion charge distribution is known. Integrating a current trace from the charge collector measurements results in the energy of the ion multiplied by the average charge state of the ion. The charge distribution $Z_i(\theta) \Delta\Omega$ averaged over a solid angle for current traces in different directions can be obtained self-consistently by comparing the charge collector data and the calorimeter data

$$Z_i(\theta) \Delta\Omega = \frac{\frac{1}{2} m_i \int i(t) dt}{E_{cal}(\theta)} \quad (6.8)$$

where m_i is the mass of an ion. Once the charge state distributions are determined, the total ablated mass is simply

$$M = \frac{m_i}{e} \iint \frac{i(t)}{Z_i(\theta) \Delta\Omega} dt d\Omega \quad (6.9)$$

$$p_{\perp} = \frac{m_i}{e} \iint \frac{i(t) v(t) \cos(\theta)}{Z_i(\theta) \Delta\Omega} dt d\Omega \quad (6.10)$$

6.1.c Procedure 3

The charge state as a function of v_i can be obtained from charge analyzer data (see chapter 4). The total ablated mass and momentum normal to the target surface are given by

$$M = \frac{m_i}{e} \iint \frac{i(t)}{Z_i(v(t)) \Delta\Omega} dt d\Omega \quad (6.11)$$

$$p_{\perp} = \frac{m_i}{e} \iint \frac{i(t) v(t) \cos(\theta)}{Z_i(v(t)) \Delta\Omega} dt d\Omega \quad (6.12)$$

6.2 Analysis of Experimental Data

Three different focus conditions (f_0 , f_1 , f_2) are used in this study to provide a wide range of laser intensities on target. The equivalent 90 % energy containing diameters for these three focus conditions are 74 μm , 170 μm and 340 μm respectively. Arrays of plasma calorimeters and charge collectors (see figure 4.1) are used to map the angular distribution of energy and velocity of the expanding plasma.

Three different procedures described in the previous section are used to deduce mass ablation rate and ablation pressure from the data. Each method highlights certain aspects of the ion data and a comparison of the results lead to a better understanding of the mechanism involved.

In procedure 1, the peak ion velocity v_p is used to estimate the average velocity $\langle v \rangle$. Unless the observed ion distribution can be described by a delta function there will be some error. The error will be small if the observed distributions consist of a sharp single peak as in the case of steady state ablation⁴⁰. As described in chapter four the ion distribution observed in this study are very complex. In any case this procedure is generally biased towards the fast ions. The ion current is given by

$$i = n_i Z_i v_i \quad (6.18)$$

where n_i , Z_i and v_i are the ion density, charge state and velocity respectively. Z_i is an increasing function of v_i and v_i is larger for faster ions, which implies that v_p in general is greater than $\langle v \rangle$. In addition, this procedure completely ignores the contribution of the slow ions. The estimated values of M and $p_\perp$ by using procedure 1 thus can be considered to correspond to the hotter plasma ablation occurring from within the laser focal spot region.

In procedure 2, the contribution of the slow ions is somewhat accounted for. But it is still biased towards the faster ions because an average charge state $Z_i(\theta)$ is used. The charge of lower energy ions is smaller than $Z_i(\theta)$ and the evaluation of M and $p_\perp$ is carried out by dividing $i(t)$ by $Z_i(\theta)$ over the whole ion current pulse.

In procedure 3, an explicit $Z_i(v)$ dependence is used, the contribution of low energy ions is fully included in this procedure. The problem is now that the measured ion traces include ions produced by the laser driven ablation as well as the ions produced by the residual ablation of the target long after the laser pulse is over.

6.3 Steady State Analytic Model

The ablation parameters for an one dimensional steady

state model (see chapter two) can be derived analytically. The mass ablation rate and ablation pressure for laser produced aluminum plasma are given by

Planar Flow

$$\dot{m}_a \text{ g cm}^{-1} \text{ s}^{-1} = 3.30 \times 10^5 (\bar{Z} \ln \Lambda \tau_{ns})^{-\frac{1}{4}} I_{13}^{\frac{1}{2}} \quad (6.13)$$

$$P_a \text{ Mbar} = 5.80 (\bar{Z} \ln \Lambda \tau_{ns})^{-\frac{1}{8}} I_{13}^{\frac{3}{4}} \quad (6.14)$$

Spherical Flow

$$\dot{m}_a \text{ g cm}^{-1} \text{ s}^{-1} = 1.03 \times 10^6 (\bar{Z} \ln \Lambda L_{\mu m})^{-\frac{2}{9}} I_{13}^{\frac{5}{9}} \quad (6.15)$$

$$P_a \text{ Mbar} = 10.17 (\bar{Z} \ln \Lambda L_{\mu m})^{-\frac{1}{9}} I_{13}^{\frac{7}{9}} \quad (6.16)$$

6.4 Numerical Calculations from 1D Hydro Code Medusa

The mass ablation rate and ablation pressure are calculated from the 1D hydrodynamic code Medusa by locating the stagnation point where the plasma flow velocities vanish. The mass ablated from the target as a function of time is calculated from the simulation results (see figure 6.2 (a)). The mass ablation rate is calculated from the following equation

$$\langle \dot{m}_a \rangle = \frac{m_a}{A_a \tau_a} \quad (6.17)$$

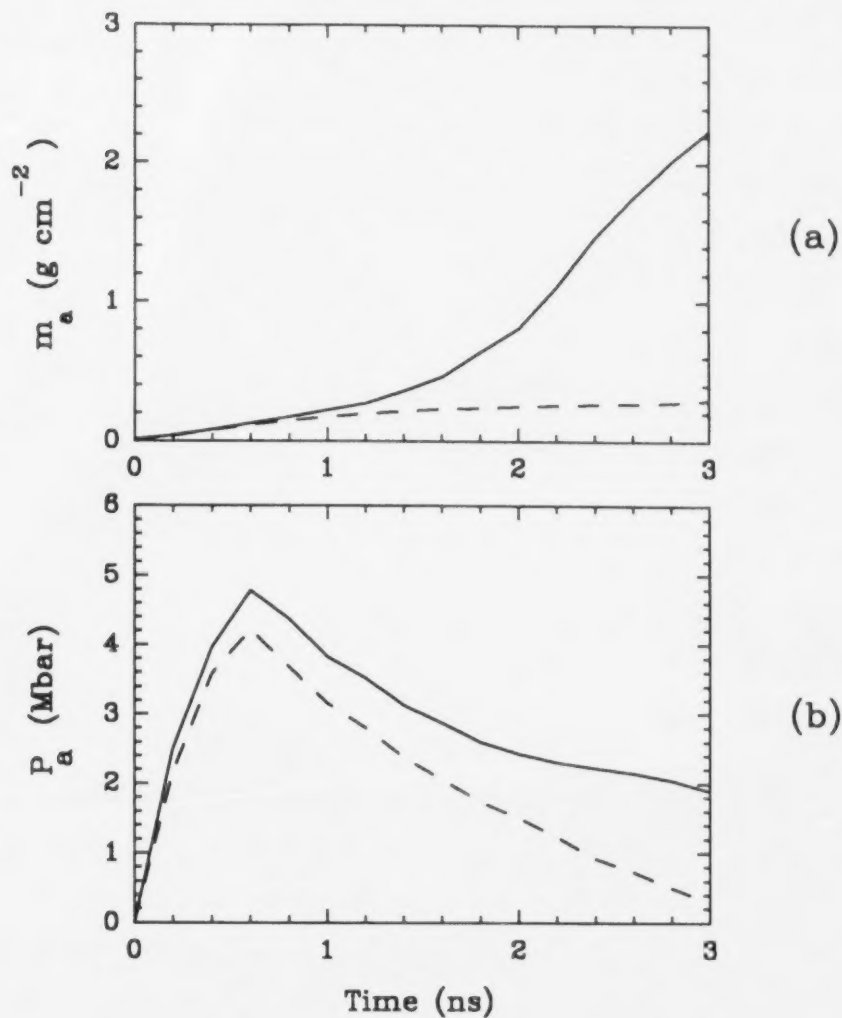


Figure 6.2 Ablation parameters of an Al plasma calculated from 1D hydro code Medusa output. The 600ps risetime triangular laser pulse has $\tau_L=1.5\text{ns}$ and $I_0=2.1\times 10^{13}\text{W/cm}^2$: (a) total mass ablated from the target as a function of time (b) ablation pressure as a function of time. The solid and dashed curves represent simulations run with and without radiation transport respectively.

The ablation time τ_a is determined in the same manner as used for obtaining the experimental results (see section 6.1). The ablated mass m_a is obtained by averaging the mass flow away from the stagnation point over the time period τ_a .

The ablation pressure is obtained by averaging the time dependent pressure (see figure 6.2 (b)) in the cell adjacent to the stagnation point over the time period τ_1 to τ_2 .

The 1D hydro code Medusa was run with and without radiation energy transport to assess its effect on the ablation parameters. The heating due to radiation energy transport ablates a large amount of mass after the laser pulse (see figure 6.2). This mass is ablated from the target with low velocity ($\ll$ ion acoustic velocity). When the contribution of the mass with a velocity less than the ion acoustic velocity is ignored, the mass ablated from the target calculated with radiation transport is about 10% higher than without it.

The ablation parameters calculated from Medusa are summarized in table 6.1. Column A gives the peak laser intensity. Columns B ,C and D, E give the mass ablation rate and ablation pressure with and without radiation transport respectively.

Table 6.1 Ablation parameters for an aluminum plasma
obtained from Medusa (1.5ns, 0.248 μ m)

A	B	C	D	E
	Without Radiation Transport		With Radiation Transport	
I_0 (W/cm ²)	$\langle \dot{m}_a \rangle$ (g/cm ² /s)	$\langle P_a \rangle$ (Mbar)	$\langle \dot{m}_a \rangle$ (g/cm ² /s)	$\langle P_a \rangle$ (Mbar)
2.1×10^{13}	7.7×10^4	2.4	2.1×10^5	3.1
1.1×10^{11}	2.1×10^4	0.07	1.0×10^5	0.10

6.5 Numerical Calculations from 2D Hydro Code Castor

Mass ablation rate and ablation pressure were calculated from the 2D hydro code Castor using the procedures mentioned in section 6.4. The time dependent results are given in figures 6.3 (a) and (b). The time averaged results are summarized in table 6.2. Columns A, B, C, D and E give the peak laser intensity, energy, pulse length, radius of the 90 % energy containing contour and average intensity of the laser respectively. The average laser intensity $\langle I_L \rangle$ for the simulation was calculated using 90 % of the energy within the 90% energy contour and the 1.5 ns pulse duration. This is the same procedure used for obtaining the experimental laser intensity on target. Columns F and G give the average mass ablation rate and ablation pressure respectively.

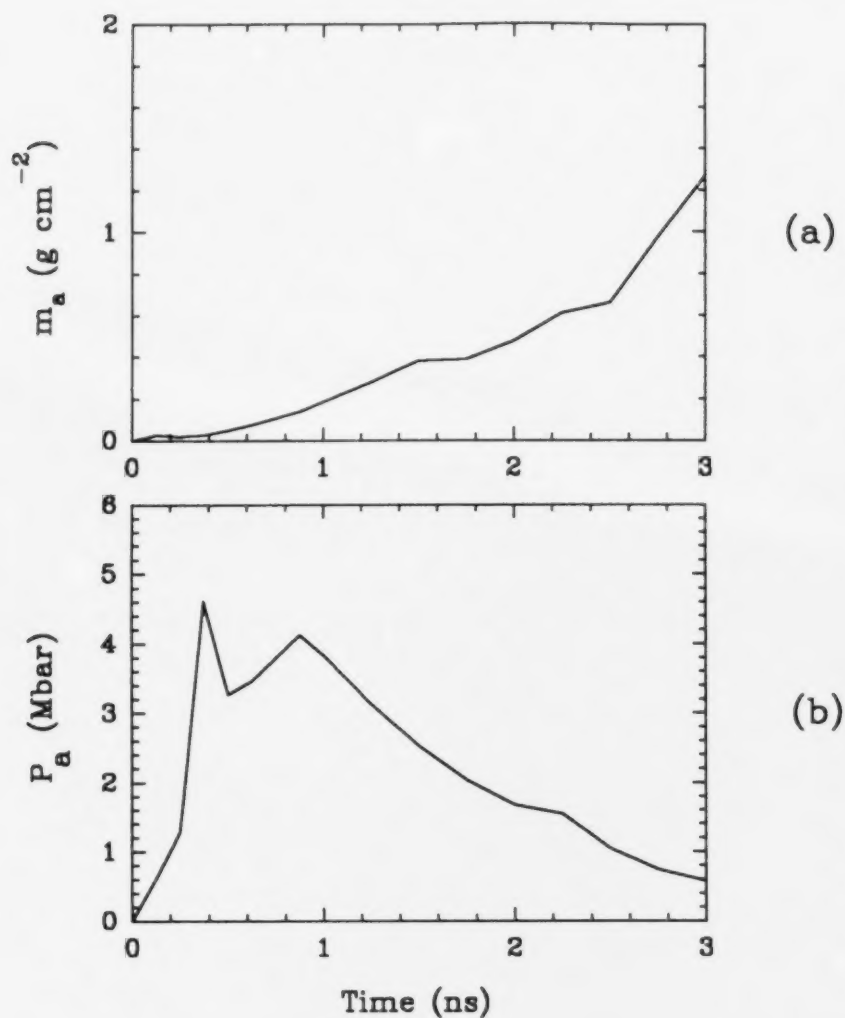


Figure 6.3 Ablation parameters of an Al plasma calculated from 2D hydro code Castor output. The 600 ps risetime triangular laser pulse has a FWHM of 1.5 ns and a peak laser intensity of $2.1 \times 10^{13} \text{ W/cm}^2$: (a) mass ablated from the target as a function of time (b) ablation pressure as a function of time.

Table 6.2 Ablation parameters from Castor for aluminum plasma produced by a 0.248 μ m laser

A	B	C	D	E	F	G
I_0 (W/cm ²)	E (J)	τ (ns)	$d_{90\%}$ (μ m)	$\langle I_L \rangle$ (W/cm ²)	$\langle \dot{m}_a \rangle$ (g/cm ² /s)	$\langle P_a \rangle$ (Mbar)
2.1×10^{13}	1.0	1.5	97	8.1×10^{12}	1.26×10^5	2.67
1.6×10^{12}	0.1	2.0	97	6.2×10^{11}	3.69×10^4	0.44
8.3×10^{11}	1.0	1.5	485	3.2×10^{11}	2.39×10^4	0.33

6.6 Summary of Ablation Parameters

The ablation parameters obtained experimentally, analytically and from simulation results are summarized in figure 6.4 and figure 6.5.

The results from 1D simulation Medusa and 2D simulation Castor are shown as hollow and solid diamond symbols respectively in figures 6.4 and 6.5.

The predictions from the 1D steady state ablation model are shown as dashed lines in figures 6.4 and 6.5. For the case of 4 mm out of focus irradiation the plasma expansion is assumed to be planar with an interaction time τ of 1.5 ns. For the cases of best focus and 2 mm out of focus irradiation, the isothermal scale length becomes larger than the focal spot radius during the period of interaction. Therefore, for these cases a spherically divergent flow with a scale length equal to the focal spot radius is assumed.

For the experimental mass ablation rates, a best fit line to the data evaluated through procedure 1 (see section 6.1.a) is shown as a solid line in figure 6.4. While this procedure will lead to a slight underestimation of mass

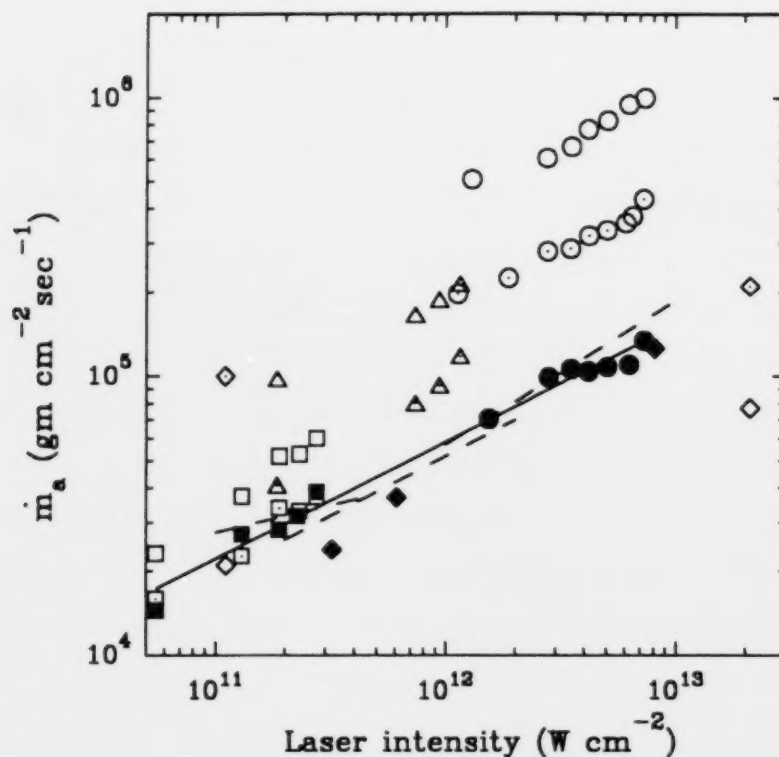


Figure 6.4 Mass ablation rate as a function of laser intensity. Circles, triangles and squares representing experimental results for three different focusing conditions f_0 , f_1 and f_2 respectively. These symbols with solid, hollow with dot and hollow representing results obtained through procedure 1, 2 and 3 respectively. The solid diamonds represent results from Castor and the hollow diamonds with and without dot on center represent results from Medusa with and without radiation transport. The dashed lines are theoretical predictions from the steady state ablation model.

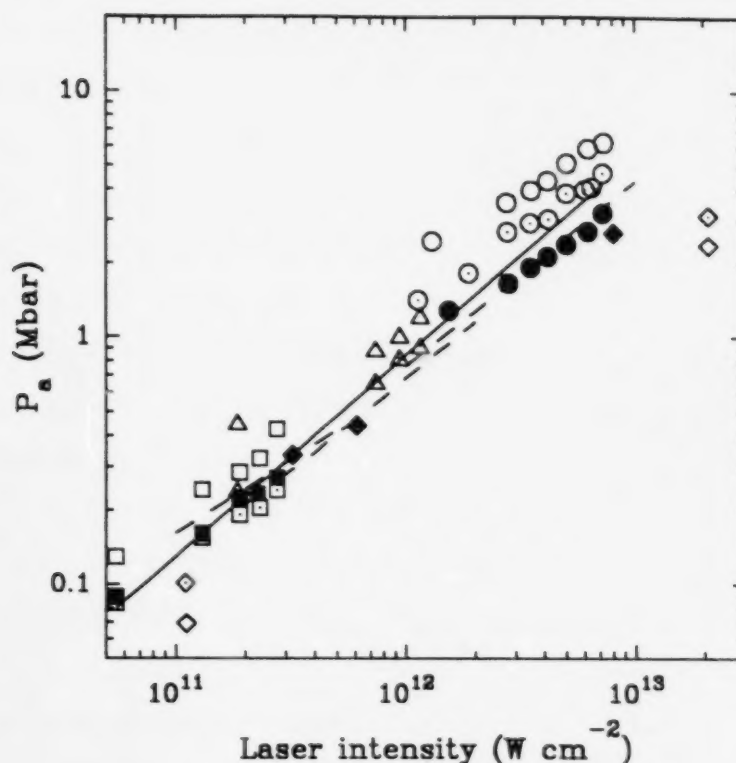


Figure 6.5 Ablation pressure as a function of laser intensity. Circles, triangles and squares representing experimental results for three different focusing conditions f_0 , f_1 and f_2 respectively. These symbols with solid, hollow with dot and hollow representing results obtained through procedure 1, 2 and 3 respectively. The solid diamonds represent results from Castor and the hollow diamonds with and without dot on center represent results from Medusa with and without radiation transport. The dashed lines are theoretical predictions from the steady state ablation model.

ablation, including data from procedures 2 and 3 (see sections 6.1.b and 6.1.c) considerably overestimates the mass ablation rate and results in a distorted scaling law. For the ablation pressure, a best fit line to the data evaluated through procedure 1 and 2 is shown as a solid line in figure 6.5. As seen in figure 6.5 the ablation pressure estimated from the two procedures differ by a factor of ~ 1.5 for the case of best focus irradiation and by a factor of ~ 1.15 for the case of 4 mm out of focus irradiation. Since these two methods represent upper and lower bounds on the true pressure, the scaling law of ablation pressure is determined by averaging the data evaluated through procedures 1 and 2. The final equations for mass ablation rate and ablation pressure are

$$\dot{m}_a = (1.5 \pm 0.1) \times 10^5 I_{13}^{0.42 \pm 0.02} \quad [\text{g cm}^{-2} \text{ sec}^{-1}] \quad (6.19)$$

and

$$P_a = (5.5 \pm 0.3) I_{13}^{0.81 \pm 0.02} \quad [\text{Mbar}] \quad (6.20)$$

6.7 Comparison with Other Data

A number of experimental studies^{46-48,64-69} have been reported on measurements of ablation parameters at UV wavelengths ($0.25\text{--}0.35\mu\text{m}$), mainly using harmonically upconverted Nd: glass laser radiation. Figure 6.5 and 6.6 give a partial list of the reported mass ablation rates and ablation pressures at UV wavelengths. Table 6.3 on page 185 summarizes the scaling laws from these data. Columns A, B, C give the wavelength, pulse length and intensity range of the lasers used respectively. Column D gives the geometry of the target with p and s representing planar and spherical respectively. Column E, F, G, H give the scaling

laws for the ablation parameters with A , α , B , β defined by the following equations

$$\dot{m}_a = A \times 10^5 I_{13}^{\alpha} \quad [\text{g cm}^{-2} \text{ sec}^{-1}] \quad (6.21)$$

and

$$P_a = B I_{13}^{\beta} \quad [\text{Mbar}] \quad (6.22)$$

The reference numbers for the publications are given in column I.

At $0.27 \mu\text{m}$ wavelength irradiation, the scaling laws are in the range of

$$\dot{m}_a = (3.1-3.3) \times 10^5 I_{13}^{(0.33-0.72)} \quad [\text{g cm}^{-2} \text{ sec}^{-1}] \quad (6.23)$$

and

$$P_a = (3.7-5.5) I_{13}^{(0.55-0.92)} \quad [\text{Mbar}] \quad (6.24)$$

At $0.35 \mu\text{m}$ wavelength irradiation, the scaling laws are in the range of

$$\dot{m}_a = (1.1-3.1) \times 10^5 I_{13}^{(0.31-0.56)} \quad [\text{g cm}^{-2} \text{ sec}^{-1}] \quad (6.25)$$

and

$$P_a = (3.7-6.7) I_{13}^{(0.66-0.84)} \quad [\text{Mbar}] \quad (6.26)$$

The intensity dependence of the scaling laws obtained in the present work ($\dot{m}_a \sim I^{0.42}$, $P_a \sim I^{0.81}$) agrees well with the range of results both at $0.27 \mu\text{m}$ and $0.35 \mu\text{m}$. The magnitude of the ablation pressure ($P_a = 5.5 \text{ Mbar}$ at 10^{13} W/cm^2) is slightly higher than previous results at $0.27 \mu\text{m}$

Table 6.3 Partial list of published scaling laws for ablation parameters at UV irradiation wavelengths

A	B	C	D	E	F	G	H	I
λ_L (μm)	τ (ns)	I_L (W/cm^2)	ge	A	α	B	β	ref
0.35	0.4	10^{13} -10^{15}	p	1.3	0.53		0.66	64
0.35	0.7	2×10^{13} 3×10^{15}	p	3.1	0.31	6.7	0.66	65, 46
0.35	1.0	1×10^{12} -2×10^{13}	p		0.53		0.77	66
0.35	0.6	2×10^{13} -10^{15}	s	1.12	0.53			67
0.35	0.6 -0.75	2×10^{13} -10^{14}	s		0.59			68
0.35	2.0	10^{12} -10^{13}	p	2.3	0.55	3.7	0.84	48
0.27	0.17	2.7×10^{14}	p	3.08	0.33			47
0.27	2.0	10^{12} -10^{13}	p	3.3	0.72	4.4	0.92	48
0.27	0.4 -1.0	5×10^{12} 10^{14}	p	1.5	0.5	3.9	0.6	69

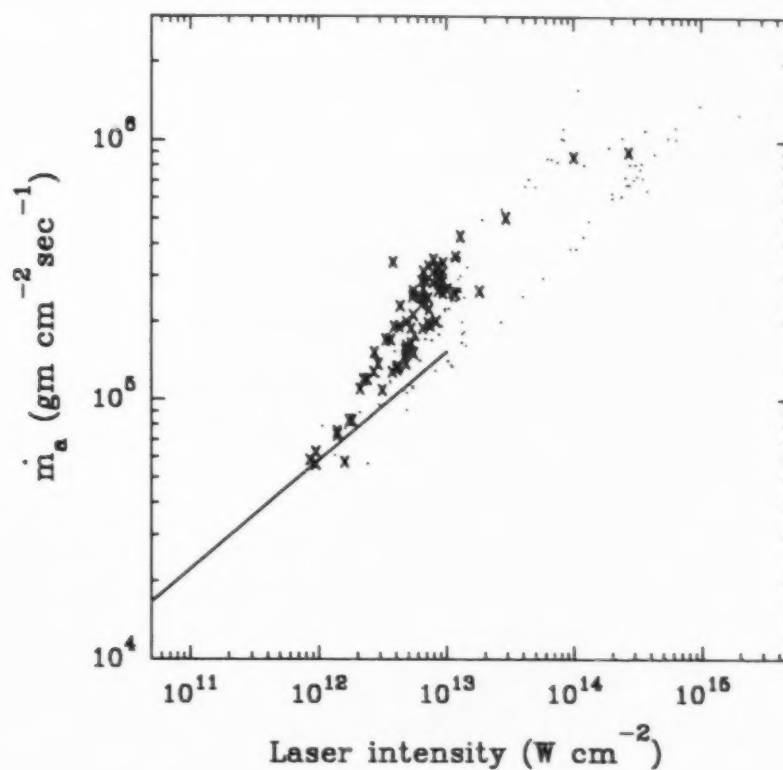


Figure 6.6 Mass ablation rate as a function of laser intensity. The dots and crosses represent published $0.35 \mu\text{m}$ and $0.27 \mu\text{m}$ results. The solid line represents the best fit of the data from this laboratory (equation 6.19).

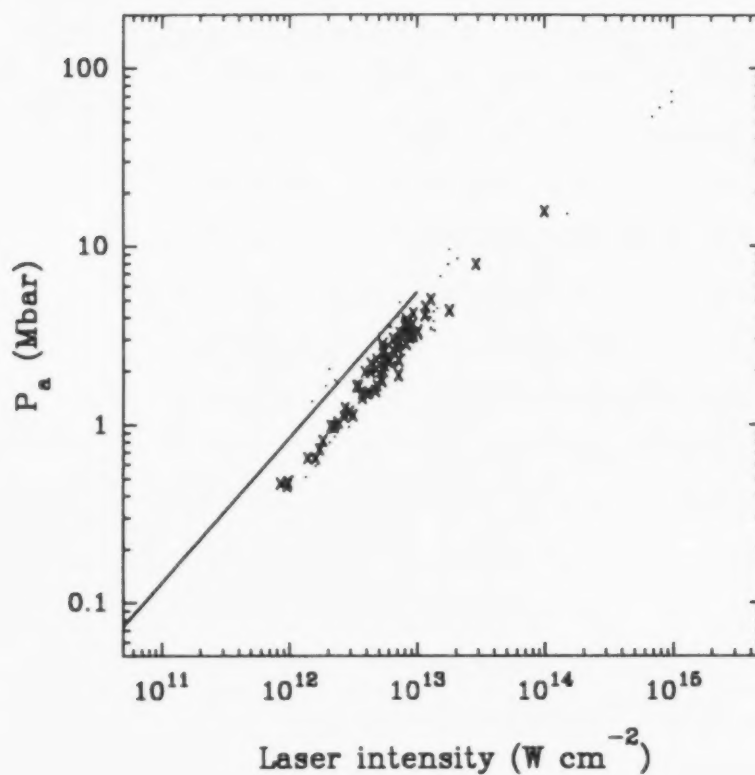


Figure 6.7 Ablation pressure as a function of laser intensity. The dots and crosses represent published $0.35 \mu\text{m}$ and $0.27 \mu\text{m}$ results. The solid line represents the best fit of the data from this laboratory (equation 6.20).

but is consistent with results at 0.35 μm . The magnitude of the mass ablation rate ($\dot{m}_a = 1.5 \times 10^5 \text{ g cm}^{-2} \text{ sec}^{-1}$ at 10^{13} W/cm^2) is lower than previous results at 0.27 μm and at the low end of the range for results extrapolated from 0.35 μm . However, as explained earlier, the mass ablation rate value derived in the present study is a lower bound and the true value is expected to be somewhat higher. Thus the agreement between the scaling laws obtained in the present study and those from the previous studies is reasonable.

6.8 Energy transport in Gold Coated Plastic Target

For low Z plasma, energy transport as measured by ablation parameters can be reasonably well described by the steady state ablation model introduced in chapter two. For high Z plasma, because of the high conversion efficiency into x-ray radiation, the steady state ablation model which does not include any radiation transport physics is not directly applicable. In order to examine the role of X-ray radiation in energy transport of high Z plasma, experiments⁷⁰ were performed using 0.248 μm KrF laser pulses of $2.0 \pm 0.2 \text{ ns}$ duration (FWHM) with energies of up to 1.2 J on 0.1 μm , 0.15 μm and 0.3 μm thick gold overcoated targets as well as plastic and gold planar targets.

Several diagnostics (see figure 6.8) were used to study the laser/plasma interaction: a luminosity time of flight spectrometer (LTOFS) viewing ionized carbon emission lines for measurements of heat penetration in layered targets; X-ray diodes and PIN diodes for XUV and soft X-ray levels; and an array of Faraday cups for the angular distribution of the plasma ion current.

Both the LTOFS and ion charge collector results indicate⁷⁰ that the laser intensity required for penetration of the high temperature ablation region through a given gold thickness is much larger than that expected from the steady state ablation model. The ion charge

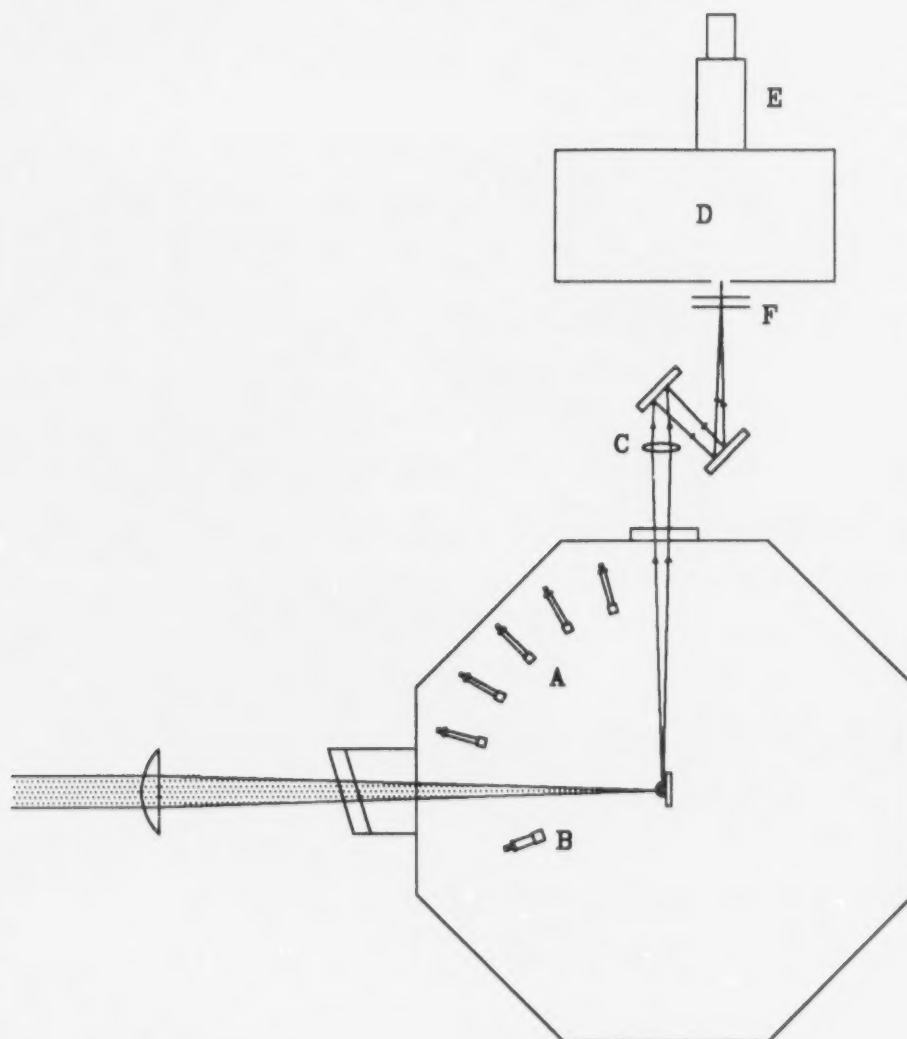


Figure 6.8 Set up for energy transport of high Z plasma experiments. A, array of Faraday cups; B, x-ray diodes. The luminosity of flight spectrometer consists of an imaging lens (c), optical filters for adjusting signal level, 0.25 m monochromator to spectrally resolve the signal and a photomultiplier to detect the signal.

collector as well as the x-ray diode results also indicate⁷⁰ the penetration of a low temperature ablation front into the plastic substrate ahead of the hot plasma burnthrough of gold overlayer. In general, the results are consistent with a two-region plasma model⁵⁰ consisting of a high temperature electron thermal transport region and a low temperature, high density radiation transport region.

6.9 Summary of Ablation Parameters and Energy Transport

Scaling laws for the mass ablation rates and ablation pressures from aluminum plasma based on ion diagnostics are obtained at $0.25 \mu\text{m}$ for intensities ranging from 10^{10} - $10^{13} \text{ W cm}^{-2}$. The results are in reasonable agreement with a 1 D steady state ablation model, numerical calculations and results from other laboratories.

Energy transport measurement for gold plasma indicates the results are consistent with a two region plasma model consisting of a high temperature electron thermal transport region and a low temperature high density radiation transport region.

CHAPTER SEVEN

DISCUSSION OF RECOMBINATION KINETICS

In this chapter experimental results are analyzed and compared to different analytic and numerical calculations. The goal is to provide a consistent description of the results leading to an understanding of the dominant physical processes involved in the expansion of a laser produced plasma to a large distance from the target surface. The main tool used in understanding this expanding plasma is the recombination kinetic code described in section 2.4.

The evolution of this expanding plasma can be divided into three main regimes. The first regime is the initial planar expansion which lasts approximately 1 ns from the beginning of the laser pulse. This is followed by the transition from planar to spherical expansion which marks the second regime. Most of the internal energy in this spherically expanding plasma will be converted into kinetic energy of expansion. When this happens the plasma can be described as undergoing an adiabatic spherical expansion which is the third regime. The hydrodynamics of the initial planar expansion and its subsequent transition to spherical expansion were discussed in section 5.5 and will not be repeated here. However, the time dependent recombination kinetics in the first two regimes is discussed in the following section. This is followed by a discussion of the recombination kinetics in the adiabatic spherical expansion phase. Subsequently a detailed comparison of simulation results with the experimental results and mechanisms contributing to the two peaked behavior observed in the aluminum ion traces are discussed.

7.1 Recombination Kinetics during Initial Heating and Expansion

The evaluation of the charge states of a high density expanding plasma is important for the interpretation of experimental results from the ion and x-ray diagnostics.

The effects of recombination kinetics on the charge state distribution is studied by comparing the charge distribution calculated from three different methods. In the first method, the time dependent ionization and recombination rate equations are solved in Lagrangian coordinates by using the density and temperature data of the plasma from the 2D hydro code Castor. In the second method the time derivative of the rate equations are equated to zero, thus allowing the ionization and recombination processes to reach steady state instantaneously. This steady state solution is obtained by using the density and temperature data from Castor. Any deviation of the results of these two calculations is a direct measure of the time dependent effects. In the third method the average value of charge state $\langle Z \rangle$ from Castor simulation is used to compare with results from the previous two methods. The $\langle Z \rangle$ from Castor is extracted from equation of state tables using a formula due to Cranfill⁷¹ in which the electrons are treated as a free semi-degenerate gas. The $\langle Z \rangle$ is given by the following empirical formula

$$\langle Z \rangle = \frac{m_i P_e}{k T_e} [1 + 0.88388 \alpha + 0.11310 \alpha^2]^{\frac{1}{5}} \quad (7.1)$$

where m_i , P_e , T_e and k are the mass of ion, electron pressure, electron temperature and Boltzmann constant and α is given by

$$\alpha = \frac{1}{2} \left(\frac{2 \pi m_e}{h^2} \right)^{\frac{3}{2}} P_e (k T_e)^{\frac{5}{2}} \quad (7.2)$$

where h is Planck's constant. The result calculated from this formula is consistent with that from Sesame equation of state table.

Figure 7.1 gives the density, temperature and $\langle Z \rangle$ as a function of time for a plasma element for the three different methods. The expanding aluminum plasma is produced by a 1.5 ns, 0.25 μm laser pulse with a peak intensity of $2 \times 10^{13} \text{ W/cm}^2$. Both the $\langle Z \rangle$ from Castor and from the steady state solution of the rate equations follow the temperature profile very closely. For these two cases as the temperature rises the charge state instantaneously reaches a He-like Al^{11+} state and the charge state then recombines to a Al^{7+} state shortly after a sharp decrease of temperature. In contrast to the steady state case, the $\langle Z \rangle$ of the time dependent case reaches the He-like charge state with a time delay and the distribution freezes even after the temperature has dropped to a low value.

A simple explanation of this result can be given in terms of the relaxation time for time dependent processes. The relaxation time τ_i of an atom staying in a particular charge state i is given by the following equation

$$\tau_i = \frac{1}{n_e (\alpha_{i+1} + S_{i-1} + \alpha_i + S_i)} \quad (7.3)$$

where n_e , α and S are the electron density, recombination rate and ionization rate respectively. In figure 7.2, the relaxation time of Al^{11+} ion is plotted as a function of density for different values of temperature. For $n_e \approx 10^{21} \text{ cm}^{-3}$ and $T_e \approx 400 \text{ eV}$ for the case presented here the characteristic ionization time should be approximately 100 ps in agreement with what is observed in figure 7.1. The

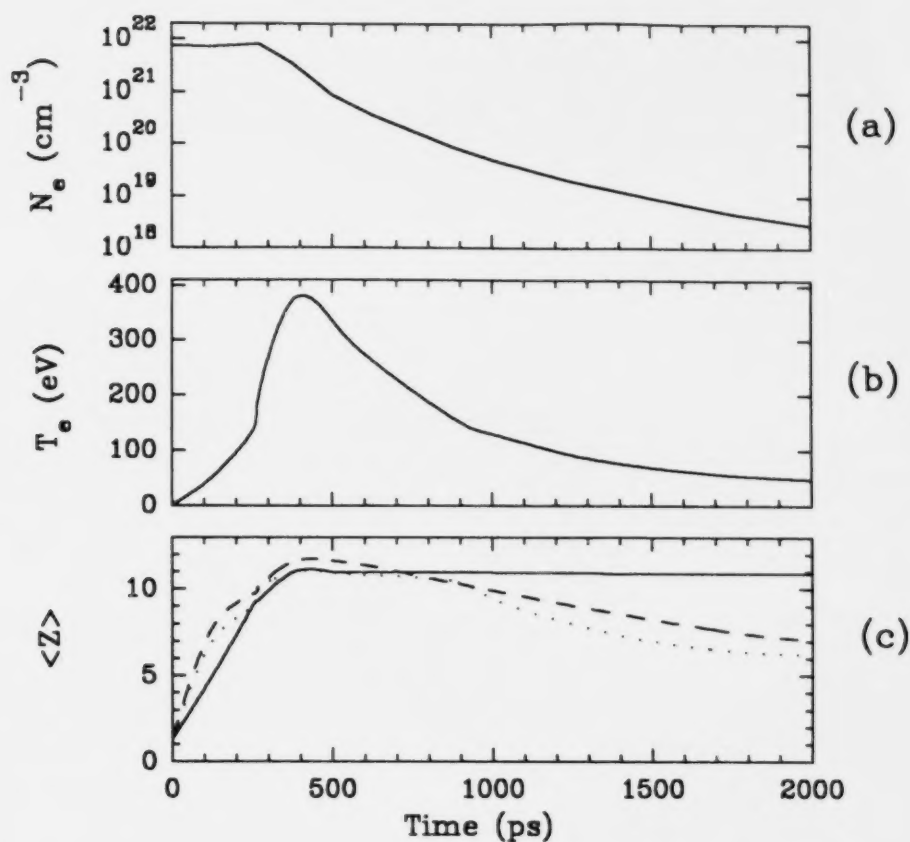


Figure 7.1 Time evolution of electron density (a), electron temperature (b) and average Z (c) for an Al plasma element with $I_{pk} = 2 \times 10^{13} \text{ W cm}^{-2}$, $\tau_L = 1.5 \text{ ns}$ and $\lambda_L = 0.25 \mu\text{m}$. In graph (c) the time dependent calculation is represented by a solid curve, the steady state calculation by a dotted curve and result from hydro code Castor by a dashed curve.

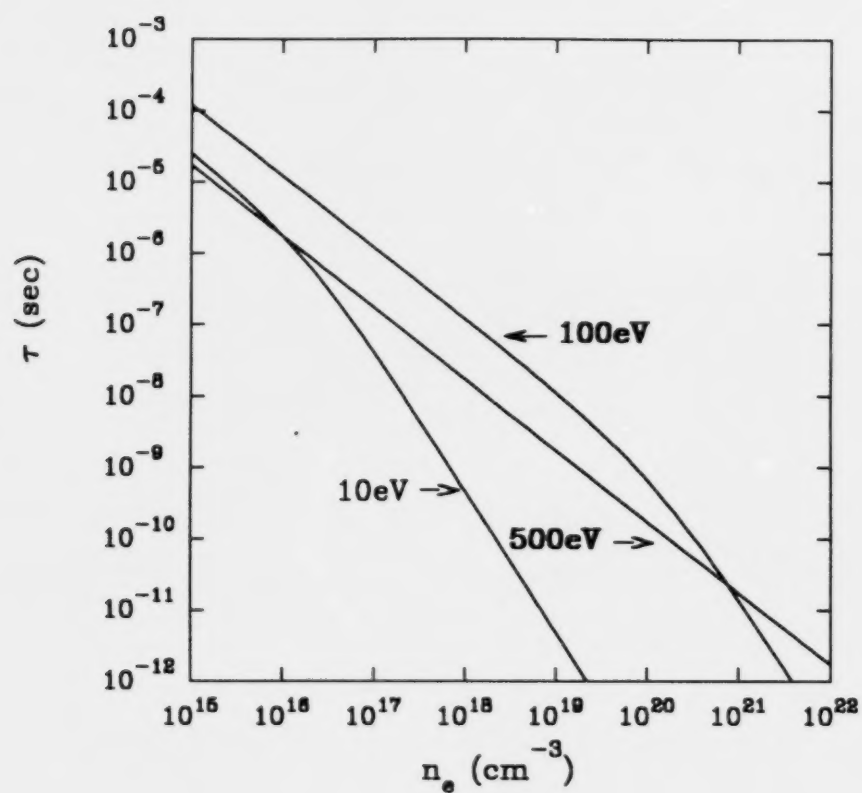


Figure 7.2 Relaxation time of Al^{11+} ion as a function of electron density.

freezing of the charge state is due to the rapid decrease of n_e as the plasma element moves away from the target surface which results in a significant increase in τ .

The freezing of the charge states are important in the determination of temperature using x-ray diagnostics. The ratio of the emitted x-ray lines $H\alpha/He\alpha$ is often used to determine the temperature of a plasma^{72,73}. Kieffer⁷³ has measured the spatially resolved temperature at distances between 20 μm to 150 μm from the target surface of a laser produced aluminum plasma using the ratio of $H\alpha/He\alpha$ line emission and the x-ray continuum emission. The x-ray continuum measurements indicate beyond 50 μm the temperature is decreasing as distance is increased from the target surface. However, the $H\alpha/He\alpha$ measurements indicate the temperature is constant up to the a distance 150 μm from the target surface. The two contradict results can be explained by the fact that a steady state model which has no freezing of ionization state is used in calculating the temperature from the $H\alpha/He\alpha$ line ratio measurement.

In addition, the freezing of the charge states are important in the ion diagnostics. A model that does not contain the effect of freezing of ionization state will underestimate the initial average charge state of the plasma and results in a lower average charge state of the plasma at a large distance from the target surface.

7.2 Recombination Kinetics of Adiabatic Spherical Expansion

As the plasma expands away from the target surface most of its internal energy will convert into kinetic energy of expansion and its expansion velocity will approach an asymptotic value.

In figure 7.3 a plasma element originating close to the target surface is tracked using its Lagrangian coordinate. This plasma element approaches an asymptotic velocity of $\approx$

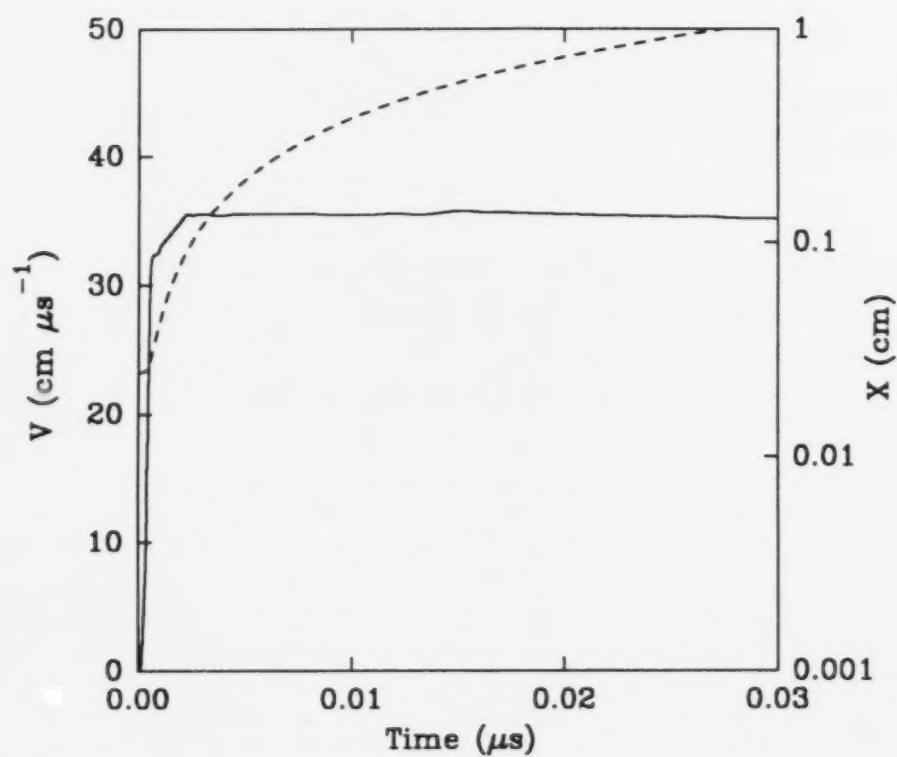


Figure 7.3 Time evolution of distance X from target surface represented by dotted curve and outward velocity v represented by solid curve for an Al plasma element with $I_{pk} = 2 \times 10^{13} \text{ W cm}^{-2}$, $\tau_L = 1.5 \text{ ns}$ and $\lambda_L = 0.25 \mu m$.

$3.5 \times 10^7 \text{ cm s}^{-1}$ at a distance $\approx 5 \text{ mm}$ away from the target surface after an expansion time of $\approx 15 \text{ ns}$. A plasma element whose expansion velocity approaches an asymptotic value can be simply described as having an adiabatic expansion.

As the plasma element expands its temperature drops and the recombination process dominates over the ionization processes. A comparison of recombination rate coefficients indicates the three body recombination dominates the radiative recombination. From equation (2.18) the three body recombination coefficients $\alpha_{3B} \sim n_e T_e^{-9/2}$. If a plasma element expands adiabatically and its equation of state obeys an ideal gas law. The temperature of the plasma element is proportional to $n_e^{2/3}$. It follows $\alpha_{3B} \sim n_e^{-2}$. Since the relaxation time τ_{3B} for the three body process is given by

$$\tau_{3B} = \frac{1}{\alpha_{3B} n_e} \quad (7.4)$$

Thus $\tau_{3B} \sim n_e$ and the rapid decrease of density of the plasma element implies τ_{3B} will become sufficiently short compared with the expansion time that the charge state distribution should reach an equilibrium state. If the charge state distribution of a plasma element is in an equilibrium state, no highly ionized charge state ions should be observed at a large distance from the target surface because of the low temperature in contrast to experimental evidence. The observation of highly ionized charge state ions at large distance from the target surface is due to the fact that during recombination a certain fraction of the ionization energy is return to the plasma as heat. This slows the decrease of temperature as a function of expansion time from $T_e \sim t^{-2}$ expected for a plasma element undergoing an adiabatic expansion to $T_e \sim t^{-1}$. This cause a reduction of recombination rate and

results in a freezing of charge state distribution as observed in the simulation and as was noted previously^{10,44,74}.

In order to confirm the assumption of adiabatic expansion is valid, two methods of simulating the ion current are used here. In the first method the time dependent rate equations and the fluid equations are solved from the first boundary close to the target surface, where the time dependent boundary conditions from Castor simulation are read in, to the second boundary located 1 m from the target surface. In the second method the time dependent rate equations are integrated from the first boundary to the secondary boundary located 1 cm from the target surface. From 1 cm to 1 m from the target surface the fluid equation are dropped and the rate equations are integrated by assuming an adiabatic spherical expansion (see section 2.5). The two ion currents simulated from these two method only differ slightly from each other which confirms the validity of assuming adiabatic spherical expansion after the expansion velocity of a plasma element reaches its asymptotic value.

7.3 Sensitivity of Rate Coefficients and Energy Return to Plasma

The various processes that are included in the recombination kinetics calculation are listed below.

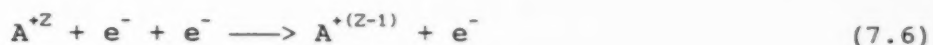
The impact ionization is the process where an electron is knocked off from an ion due to collision with another electron. The process is given by



Experimental results for impact ionization rates are generally very limited. Several theoretical methods are available for calculating the rates. They are (1) the

semi-empirical Lotz formula⁷⁵, (2) the distorted-wave exchange approximation⁷⁶, (3) the scaled Coulomb Born approximation⁷⁶ and (4) the exchange classical impact parameter method⁷⁶.

Three body recombination is the inverse process of the electron impact ionization and is given by



Its rate can be obtained from the ionization rate by the method of detailed balance in thermodynamic equilibrium. The various ionization and three-body recombination rates are generally in agreement within a factor of approximately two.

While the three body recombination is important for high density plasma, in low density the plasma the dominant recombination processes are radiative and dielectronic recombination. The radiative recombination is the process where an electron is captured by an ion and the excess binding energy is emitted as a photon. The process is given by



This recombination process is the most important one for the fully ionized, hydrogen-like and helium-like ions. Its rate is usually calculated by using the Kramer formula⁷⁶ and generally believed to be quite accurate.

The binding energy of an ion that is recombined from charge state i to a lower charge state $i-1$ is defined as

$$E_i^B = \chi_i - \chi_{i-1} \quad (7.8)$$

where χ_i is the energy required to ionize an atom from neutral to charge charge i . During the recombination process a portion of this energy is return to the plasma

while the remaining is allowed to escape as line radiation. The energy return to the plasma through recombination is usually derived⁹ by assuming the plasma is optically thin to the x-ray line radiation. The inclusion of opacity of the plasma should increase the amount of returned energy. In any case its upper limit is the binding energy. Figure 7.4 gives the fraction of the binding energy returned to the plasma for an expanding aluminum plasma element. The plasma is produced at the following irradiation conditions: $I_{pk} = 2 \times 10^{13} \text{ W cm}^{-2}$, $\tau_L = 1.5 \text{ ns}$ and $\lambda_L = 0.25 \text{ }\mu\text{m}$. The fraction of binding energy returned to the plasma is defined as

$$F^{ER} = \frac{\sum_i f_i E_i^*}{\sum_i f_i E_i^B} \quad (7.9)$$

where E^* is given by equation (2.19). F^{ER} is ≈ 0.2 in early time of the expansion and subsequently decreases to ≈ 0.001 in late time of the expansion. Latyshev⁷⁷ suggested a different E^* can be calculated if a different rate for the radiation transition of the electron through the atomic energy level (see section 2.4.b.2) is used.

The various published rate coefficients and expressions for heat return to the plasma are only in agreement with each other within about a factor of two. Thus, the rate coefficients and heat return to the plasma are allowed to vary from 0.5 to 2.0 times their nominal values in order to quantify the effect of their variation on the simulated ion currents. For the case of aluminum plasma with irradiation condition $I_{pk} = 2 \times 10^{13} \text{ W cm}^{-2}$, $\tau_L = 1.5 \text{ ns}$ and $\lambda_L = 0.25 \text{ }\mu\text{m}$, a factor of two variation of the ionization rate coefficient and the radiative recombination rate coefficients produce no significant change in the simulated ion current, a factor of two variation of the three body recombination rate coefficients and the heat return to the

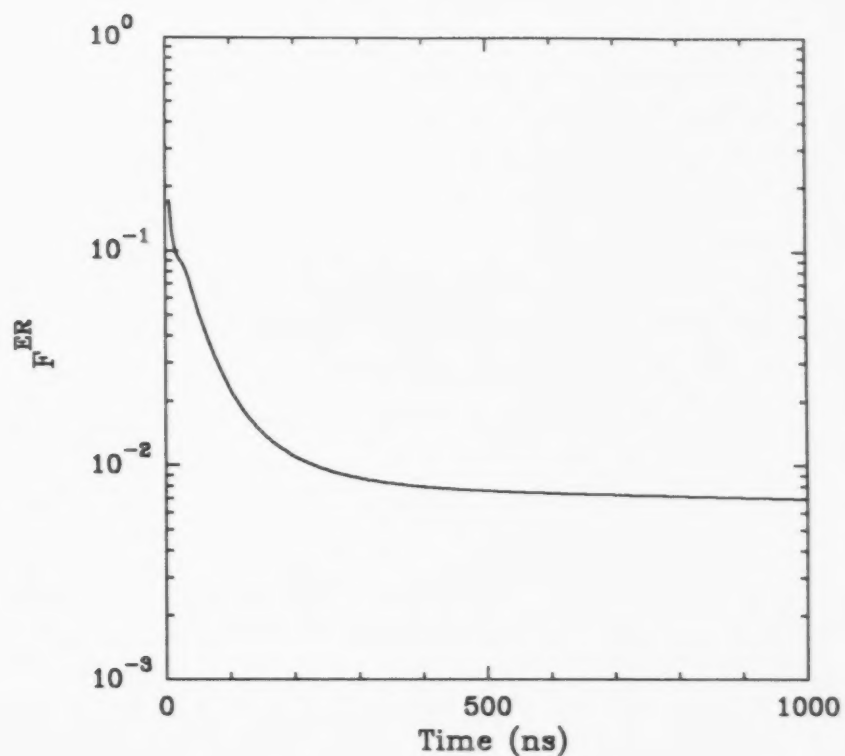
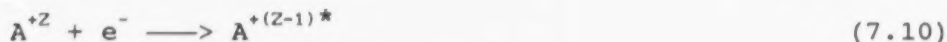


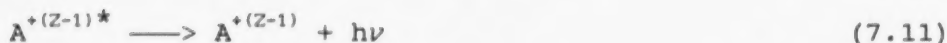
Figure 7.4 Fraction of the binding energy returned to the plasma for an expanding aluminum plasma element produced at the following irradiation conditions: $I_{pk} = 2 \times 10^{13} \text{ W cm}^{-2}$, $\tau_L = 1.5 \text{ ns}$ and $\lambda_L = 0.25 \text{ } \mu\text{m}$. The fraction of binding energy returned to the plasma, F^{ER} is defined in the text.

plasma results in a 9 % and 24 % change of the average Z value of the simulated ion current respectively. This indicates that the heat return to the plasma is the dominant factor in this recombination kinetic calculation.

One process that is not included in the simulation is dielectronic recombination. It is a process in which an electron is resonantly captured by an ion in a level with large principal quantum number and the excess binding energy is used to excite another bound electron in the ion to a excited state. The excited ion is stabilized by emitting a photon. The processes are given by



and



Dielectronic recombination operates in the same region where the radiative recombination is important. Its rate is directly proportional to the electron density and thus the rate coefficient can be added to that of radiative recombination. The calculation of the rate coefficients involves summing over all intermediate excited states using oscillator strengths for all transitions. Since this difficult calculation requires a large amount of atomic data, the use of different sets of atomic data can result in the dielectronic recombination rates differing by as much as a factor of four. However, the fact that the radiative recombination is not a dominant process in our calculation indicates the omission of the dielectronic recombination will not lead to a large error.

7.4 Peak Velocities of Ion currents

A comparison of simulated ion currents calculated with

the recombination kinetics and without recombination kinetics (constant Z model) indicates firstly the signal from the constant Z model is higher and secondly the peak velocity from the constant Z model is slightly slower. For the case of aluminum plasma with irradiation condition $I_{pk} = 2 \times 10^{13} \text{ W cm}^{-2}$, $\tau_L = 1.5 \text{ ns}$ and $\lambda_L = 0.25 \mu\text{m}$ the shift of peak ion velocity is $\approx 10\%$. This can be explained simply by noting that within a charge state level, the slower ions experience more recombination than the faster ions⁴⁴. This results in a shift of ion peak to higher velocity with increasing distance from the target surface.

7.5 Comparison with Experimental Results

The individual charge state ion traces were observed using a charge analyzer and the combined ion trace was observed using a Faraday cup. The ion traces for a carbon and an aluminum plasma at $I_L \approx 8 \times 10^{12} \text{ W/cm}^2$, $\tau_L \approx 2 \text{ ns}$ and $\lambda_L \approx 0.25 \mu\text{m}$ are shown in figures 7.5 and 7.6. The sum of individual charge states reproduced the Faraday cup trace that was observed at similar laser conditions. The combined charge state trace of the aluminum plasma show a two peak behavior. This two peak behavior will be discussed later in section 7.6. No two peak behavior was observed in the case of the carbon plasma.

Figure 7.7 and 7.8 show the simulation results of carbon and aluminum plasma respectively for peak laser intensity of $2 \times 10^{13} \text{ W cm}^{-2}$ which is equivalent to an average laser intensity of $8.7 \times 10^{12} \text{ W cm}^{-2}$ if the criteria that defines the experimental laser intensity is used. In general the overall behavior of the ion current as a function of time from the simulation agrees quite well with the experimental measurements. However, comparing these results in detail, there are discrepancies in the charge state spectrum. This discrepancy is particularly pronounced in the late time charge state spectrum where the

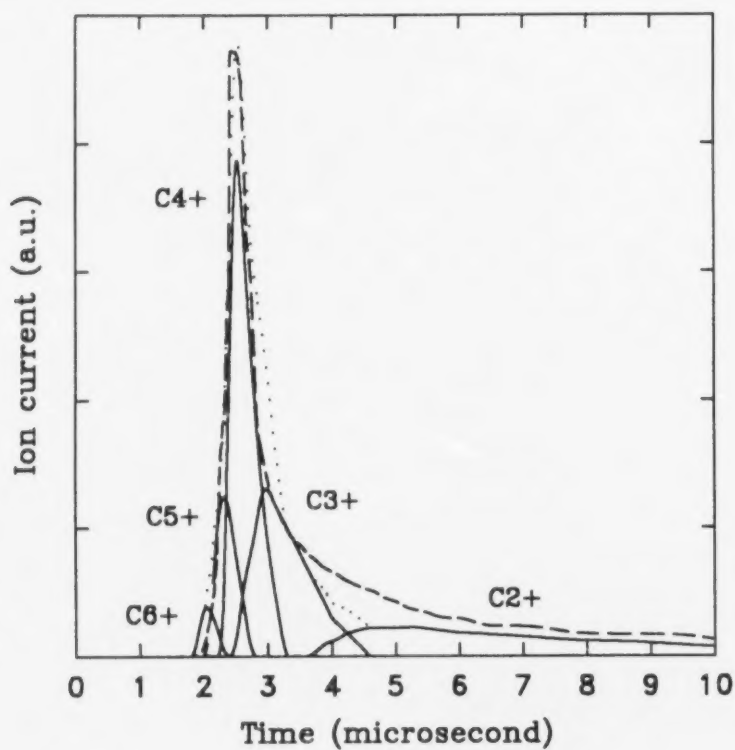


Figure 7.5 Charge state traces observed normal to the target surface for a carbon plasma at a laser intensity of $7.3 \times 10^{12} \text{ W/cm}^2$. The dotted line is the sum of charge state traces and the dashed line is the Faraday cup trace under similar experimental conditions.

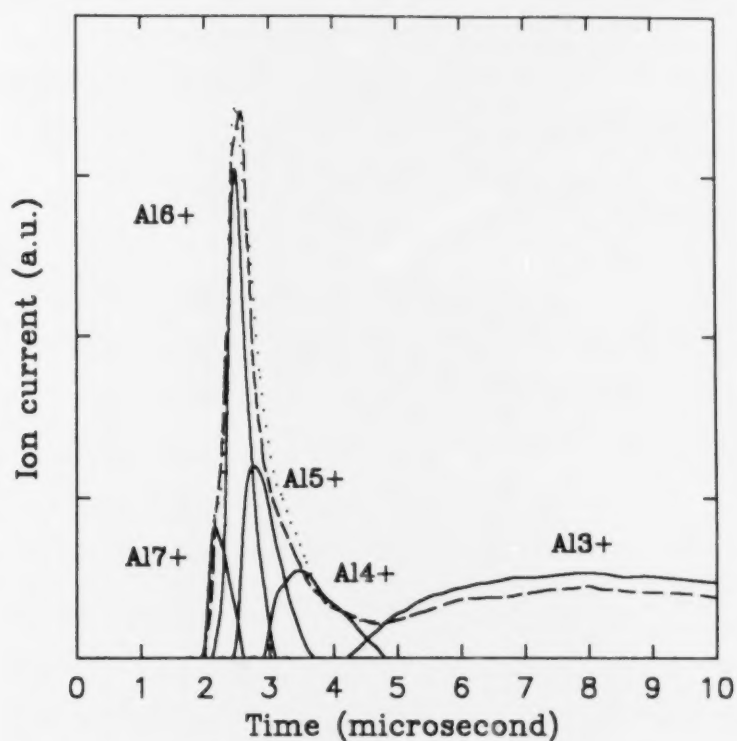


Figure 7.6 Charge state traces observed normal to the target surface for a aluminum plasma at a laser intensity of $7.5 \times 10^{12} \text{ W/cm}^2$. The dotted line is the sum of charge state traces and the dashed line is the Faraday cup trace under similar experimental conditions.

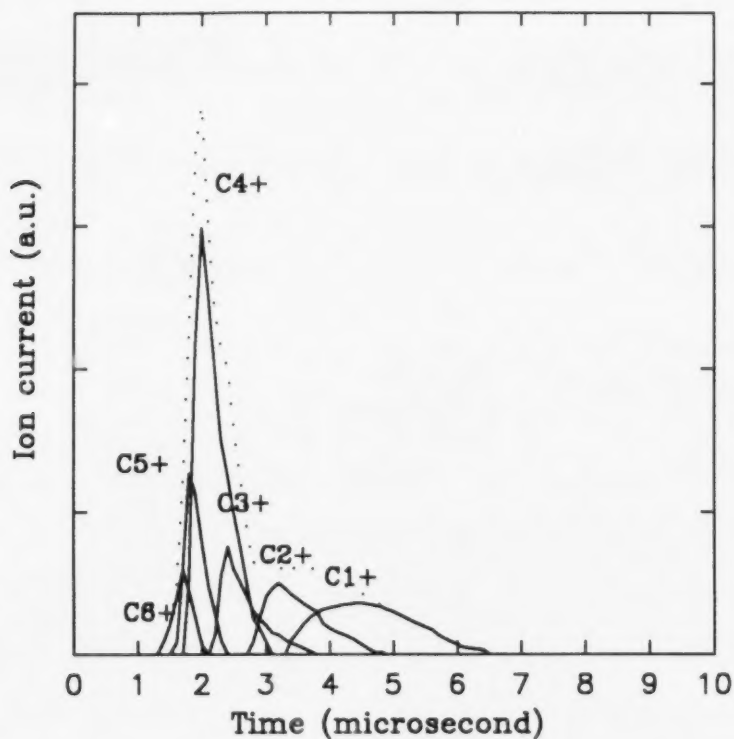


Figure 7.7 Charge state traces of C plasma from recombination kinetics simulation at a peak laser intensity of $2 \times 10^{13} \text{ W cm}^{-2}$ corresponding to a laser intensity of $8.7 \times 10^{12} \text{ W cm}^{-2}$ calculated using 90% of the energy and the area of the 90% energy contour. The dotted line is the sum of charge state traces.

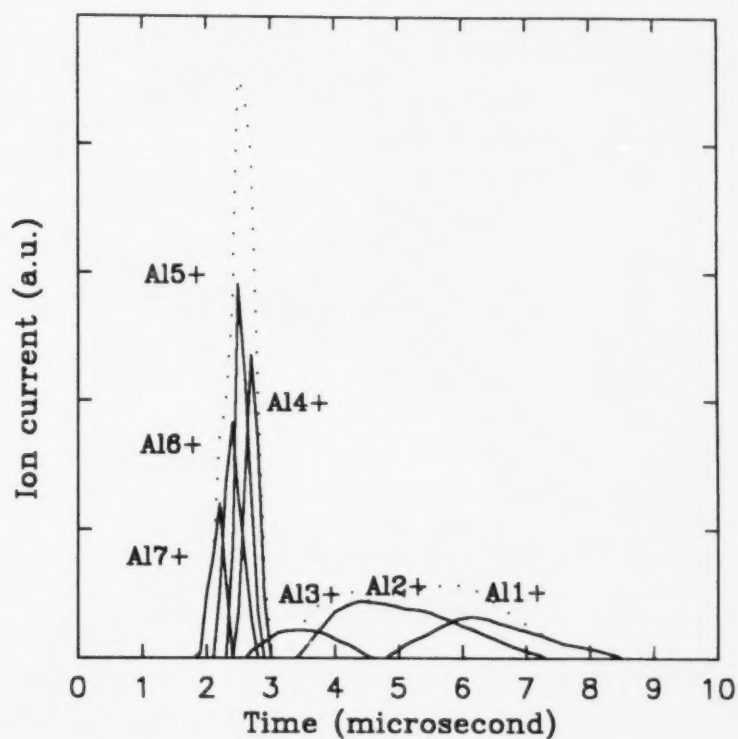


Figure 7.8 Charge state traces of Al plasma from recombination kinetics simulation at a peak laser intensity of $2 \times 10^{13} \text{ W cm}^{-2}$ corresponding to a laser intensity of $8.7 \times 10^{12} \text{ W cm}^{-2}$ calculated using 90% of the energy and the area of the 90% energy contour. The dotted line is the sum of charge state traces.

experimental results give higher charge states than the simulation results. The observed discrepancy between experimental and simulation results may be due in part to the fact that a simplified 1 D calculation is used to approximate a 3 D expansion resulting in not very accurate plasma conditions.

Figure 7.9 (a) and (b) show the average Z as a function of time for aluminum and carbon plasma respectively from experimental data (dots) and its best linear fit (solid lines) and from simulation results (dashed lines). The agreement for the initial period from 2 to 3 μs is reasonable. However for time $> 3 \mu\text{s}$ the experimental results is somewhat higher. This is probably due to the fact that the simulation laser pulse is completely turned off at the end of 3 ns leading to rapid temperature decrease in the corona. In contrast the experimental laser pulse has a longer tail which prevents the late time plasma from cooling so rapidly. In addition, the simulated ion current for time $> 4 \mu\text{s}$ originates from plasma produced after 6 ns from the onset of the laser pulse and is quite sensitive to the extrapolation method that is used for the initial plasma conditions after 6 ns.

7.6 Ion Current at Non-normal Angles of incidence

The ion currents at different observation angles are measured and reported in chapter four. The average charge state of the ions is approximately linearly proportional to the velocity of the ions without any apparent dependence on observation angle. The peak ion velocity as well as its average charge state decreases with increasing observation angle. In addition, the width of the ion current pulse increases with increasing observation angle. This is to be expected since the ion current at different observation angles originates from different regions of the initial plasma during the laser heating as implied by the plume

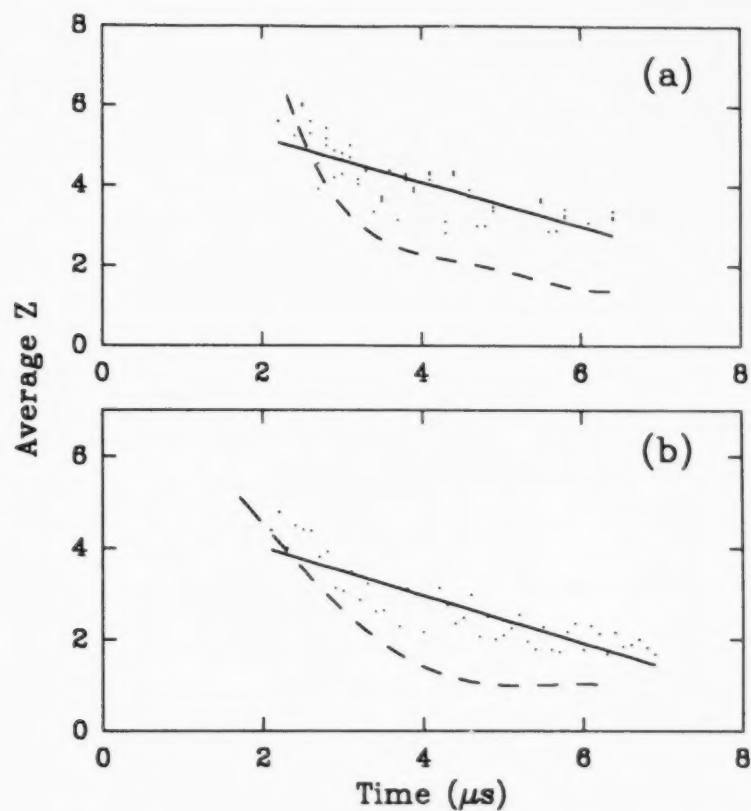


Figure 7.9 Average Z as a function of time for aluminum (a) and carbon plasma (b) respectively from experimental data (dots) and its best linear fit (solid lines) and from simulation results (dashed lines). The plasma are produced by a laser pulse of $\lambda = 0.25 \mu m$, $\tau \approx 2 ns$ and $I \approx 8 \times 10^{12} W cm^{-2}$.

model. The ion current from a larger observation angle maps the region of the initial plasma that is further away from the laser axis⁷⁸. The temperature of the initial plasma decrease radially away from the laser axis. This means the initial plasma temperature decreases for ion current at larger observation angles. This decrease in initial temperature as the observation angles are increased implies the ions will be slower as observed. The width of the ion peak was also seen to increase with larger observation angle which may be related to the larger temperature gradients off axis in the source plasma. It also affects the charge state distribution at a large distance from the target surface in two ways: (1) a cooler initial temperature results in a lower average charge state distribution and (2) slower ions results in stronger recombination at a fixed distance from the target surface as evidenced by the monotonic increase of average charge state as a function of velocity. Both these processes contribute to the decrease in average charge state for the peak ion current at increasing observation angles.

As discussed earlier in section 7.1, correct time dependent boundary conditions for the recombination kinetic code that include the effect of time dependent ionization and freezing of charge state requires the tracking of plasma elements in Lagrangian coordinates and integrating the time dependent rate equations using the corresponding density, temperature and velocity data from Castor. To obtain these boundary conditions for off axis plasma elements is a complex 2D problem and is not pursued in this thesis work. However, a simplified approach was used to simulate these ion currents at different observation angles. At the spatial boundary of the Castor simulation, 600 μm from the target surface the plasma that has a velocity vector pointing at the observation angle is assumed to end up in the ion detector at that angle. For a given time, the charge distribution is $f_{i-1} = \zeta$ and $f_i = 1$

- ζ where $i-1 < \bar{Z} < i$ and $\bar{Z}$ is the average charge state provided by Castor and ζ is derived according to the following formula

$$\bar{Z} = \zeta f_i + (\zeta - 1) f_{i-1} \quad (7.12)$$

that is, if the Castor average Z is for example 7.4, then a charge distribution of $f_7 = 0.6$ and $f_8 = 0.4$ is assumed. The simulated ion currents using the above approach have average Z much smaller than the measured values. This is probably due to the fact that the boundary charge state distribution obtained using the above method does not include the effect of freezing of charge state distribution which results in a considerable underestimation of the initial charge state distribution.

7.7 Ion Distribution for Long Pulse Irradiation

The ion current measurements for long pulse, low intensity KrF and Ruby cases are reported in chapter four. In figure 7.10 the individual charge state ion traces obtained using a charge analyzer and the combined ion trace observed using a Faraday cup are plotted. The ion traces are for a carbon plasma at a laser intensity of $\approx 1 \times 10^{11}$ W/cm² for a 25 ns, 0.694 μ m laser pulse. The sum of individual charge states reproduced the Faraday cup trace that was observed under similar laser conditions. The small discrepancy at the end of the pulse is probably caused by loss of signal due to space charge spread for very slow ions.

For the high intensity short pulse irradiation condition, the measured ion current has a sharp pulse characteristic of steady state ablation. The ion currents for the long pulse cases consist of a much wider pulse. This is to be expected since the lower laser intensities result in a lower plasma temperature yielding slower ions

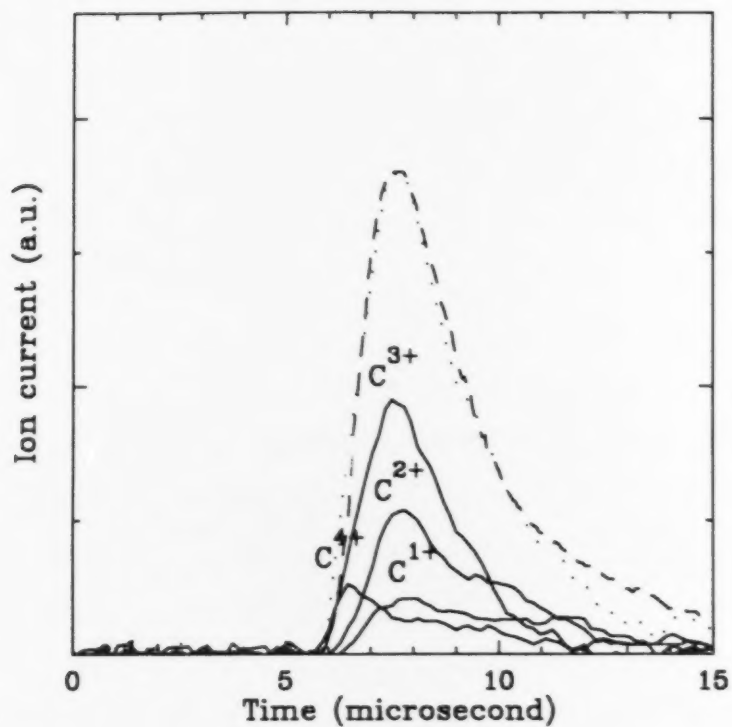


Figure 7.10 Charge state traces observed normal to the target surface for a carbon plasma at $I_L = 1 \times 10^{11} \text{ W/cm}^2$, $\tau_L = 25 \text{ ns}$ and $\lambda_L = 0.694 \text{ }\mu\text{m}$. The dotted line is the sum of charge state traces and the dashed line is the Faraday cup trace under similar experimental conditions.

for which the velocity spread is larger. A comparison of the long pulse KrF at $I_L \approx 1.7 \times 10^{11} \text{ W cm}^{-2}$ and the long pulse ruby case Ruby case at somewhat lower intensity $I_L \approx 1 \times 10^{11} \text{ W cm}^{-2}$ indicates the peak ion velocity and the average charge state of the peak ions for the ruby case is somewhat higher. This result is to be expected since according to the steady state ablation model the temperature of the plasma is $\propto (I_L \lambda_L)^{4/9}$ which implies the plasma temperature for the ruby case should be somewhat higher. In addition, the initial plasma density should be lower for the ruby case since according to the steady state ablation model the density for the sonic point of the plasma is $\propto (I_L/\lambda_L^2)^{1/3}$. Both of these factors result in weaker recombination for the ruby case. The angular distribution of the ion currents for the long pulse case share similar qualitative feature as the short pulse case and can be explained using similar arguments.

7.8 Two Peak Behavior

At high laser intensity irradiation the aluminum ion traces show a two peak behavior. It was observed that the first peak becomes increasingly dominant with increasing laser intensity. It was also observed that as the direction of observation was changed, the second peak become relatively more prominent with increasing values of observation angle. No two peak behavior was observed in the large focal spot case.

An insight into the two-peak behavior can be obtained by calculating the velocity distributions from the ion current traces (figure 4.25 (b), (d) and (f)). The velocity distributions corresponding to the three focal conditions: best focus (f0), 2 mm off-focus and 4 mm Off-focus (f2) are shown in figure 7.11. It is seen that for the high-intensity best-focus irradiation, The ion velocity distribution show a plateau region which lies between the

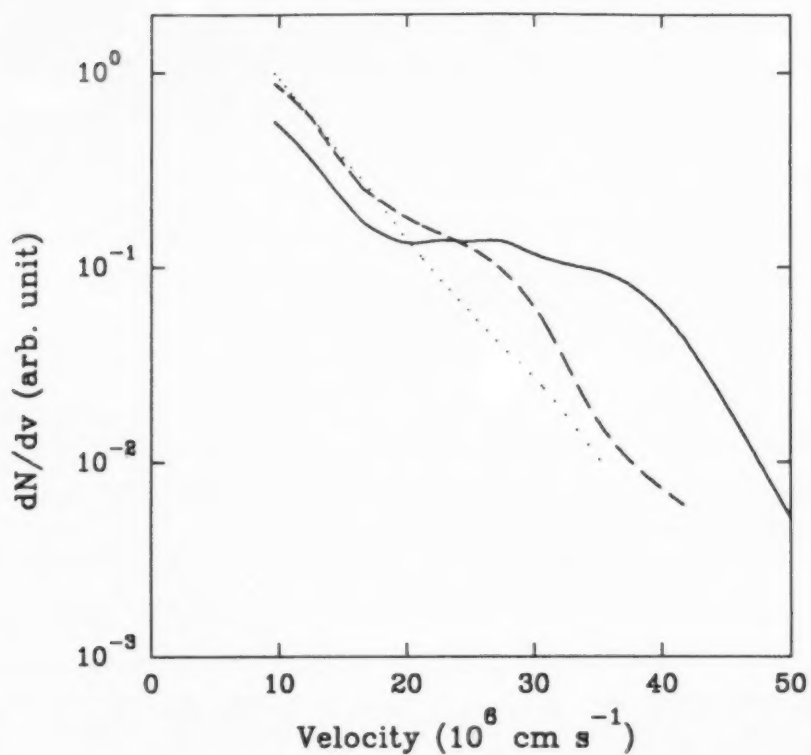


Figure 7.11 Ion velocity distributions for Al ions observed in a direction 14° to the target normal. The laser intensities (2 ns, $0.268 \mu\text{m}$) for the three curves (solid, dashed and dotted) are $7.1 \times 10^{12} \text{ W/cm}^2$ (f0), $1.1 \times 10^{12} \text{ W/cm}^2$ (f1) and $3.7 \times 10^{11} \text{ W/cm}^2$.

velocities corresponding to the two ion peaks. At lower intensities the length of the plateau region decreases. Furthermore, The velocity distribution in the region of the second peak fits a straight line on the semilog plot, which is typical of an isothermal expansion. For the low-intensity large-focal-spot irradiation, the velocity distribution does not shown any plateau region.

Estimates of the relative energy content of the ions of the two peaks were obtained by splitting the current traces into two overlapping peaks, such that their sum represents the observed ion traces. The kinetic energy and the momentum normal to the target of the ions were then computed using $Z(v)$ from chapter four. It is observed that whereas ions in the first peak are emitted preferentially closer to the target normal, ions in the second peak show a more isotropic distribution. Angular integration of the energy revealed that the ion energy in the first peak increases with laser intensity, reaching 50 % of the total for the maximum laser intensity of the experiment.

The possibility of ions in the first peak being protons due to an impurity layer of hydrocarbon on the target surface under irradiation can be ruled out. If we integrate the current in the first peak and use $Z_1 = 1$ (for protons), the estimated thickness of the layer would be $1 \mu\text{m}$ which is much more than the thickness of contamination expected for a clean surface. This was confirmed by comparing the ion current signals observed from a fresh target surface to those observed when the target surface was heated prior to irradiation. The preheat was done by firing ten shots of defocused laser beam to given an energy density of $\sim 10 \text{ J/cm}^2$, sufficient to evaporate the surface layers. No significant change in the two-peak behavior of signals was observed.

No significant parametric instabilities are expected in the present experimental conditions and no significant inhibition in thermal conduction is expected for

intensities in our experiments. Self-generated magnetic fields are not expected to have any significant effect in modifying energy transport in our high density plasma.

The second peak cannot be due to the unloading of shock heated material since only $\approx 3\%$ of the absorbed energy is contained in the shock heated material versus the $\sim 50\%$ energy observed in the second peak. Also the estimated initial temperature of the plasma leading to the second peak is $T_e \sim 50 - 100$ eV which is much greater than the few eV expected from shock heating.

The observed two peak behavior can be caused by three possible mechanisms: (1) recombination kinetics in the expanding plasma. (2) lateral heat conduction in the plasma, (3) plasma ablation by x-ray radiation transport at a lower temperature than that caused by direct laser radiation.

7.8.a Recombination Kinetics

The simulated aluminum ion trace indicates a smaller second peak than observed experimentally. However if the recombination kinetics were such that the tail of the pulse remained in Al^{3+} charge state the current in the tail would increase reproducing in part the observed second peak. However the simulated ion current tail is significantly shorter than the experimental observed tail. This indicates that recombination kinetic may be in part responsible for the occurrence of this two peak behavior. However, this is sensitive to extrapolation of the starting plasma conditions after 6 ns.

7.8.b. Lateral Heat Transport

The observed two-peak behavior may be qualitatively explained by considering lateral heat conduction in the plasma. When a laser beam irradiates a planar target, initially a local region of hot plasma is formed. Due to

the high temperature in this plasma, thermal diffusion is important. The scale length for thermal diffusion L_T is given by $\sqrt{4 D t}$, where D is the thermal diffusivity and t is the time to establish the steady flow pattern. When $L_T > R$, one would expect the heat zone to increase in radius until it reaches a dimension of the order of L_T . However, the increased effective size of the heated region results in a steady-state three-dimensional (3D) flow structure in the plasma with an effective size larger than the initial focal spot. The plasma ablated in this state would have characteristics corresponding to a reduced effective laser intensity. Under these conditions, a transition is expected from the initial high-temperature flow to a lower-temperature steady-state flow. On the other hand, if the initial focal spot size is much larger than L_T , the flow will remain essentially 1D.

The thermal diffusion length can be estimated for the typical parameters of our experiment. For a laser intensity of 10^{13} W/cm² at best focus, $T_e = 330$ eV, $n_e = 4.5 \times 10^{21}$ cm⁻³, $Z_1 = 11$, the thermal diffusion length is ≈ 100 μ m, which is much larger than the focal spot radius of 37 μ m. Thus, by considering thermal conduction, lateral energy flow would certainly be very important for best-focus ($R = 37$ μ m) and the 2-mm out-of-focus ($R = 87$ μ m) conditions and less important for the 4-mm-out-of focus ($R = 170$ μ m) condition.

It is necessary, however, to include energy dissipation due to ablation in the periphery, which would have the effect of reducing the size of the heated region. An estimate of this modified thermal conduction length can be obtained by comparing the lateral thermal conduction heat flux to the outward convective energy flow in a manner similar to that was calculated by Pert⁷⁹. By balancing the thermal heat flux KT_e/L_T over a conduction length of L_T , where K is the heat conductivity, and the convective flow of energy in the periphery of the plasma, $3 m_i n_i c_s^3$, where

m_i is the ion mass and n_i is the ion density, the conduction length is estimated to be $\sim 54 \mu\text{m}$ for an average temperature in the periphery which is half that in the central region.

The actual heat zone is expected to lie somewhere between the above two limits. Results from the 2D numerical calculation Castor indicates the scalelength is $\approx 70 \mu\text{m}$ during the laser peak and in the corona where the temperature is maximum.

Thus, the ions in the first peak may correspond to the initial hot ablating plasma still undergoing 1D flow, while those in the second peak may correspond to the late time 3D plasma flow occurring over an effectively large area.

However, under similar irradiation conditions the carbon ion current shows no signs of two peak behavior. Further reduction of focal spot radius to $10 \mu\text{m}$ by using a 20 cm focal length lens did not result in two peak behavior in carbon ion current. These observations indicate the lateral transport via electron thermal conduction cannot be the dominant process responsible for the two peak behavior in the aluminum ion current traces.

7.8.c Ablation by X-ray Radiation Transport

As discussed in chapter five, 20 % of the absorbed energy in the aluminum plasma is emitted in x-ray. Thus radiation transport in aluminum may partly be responsible for the two peak behavior in the ion current trace. Figure 7.12 gives simulated aluminum ion currents by using plasma conditions calculated by Medusa with and without radiation transport. A weak two peak behavior is observed in the ion current simulated from plasma condition with radiation transport turned on. The one dimensional planar nature of the Medusa simulation implies that heat dissipation via convection is very slow. This will result in an unrealistically high coronal temperature that produces an

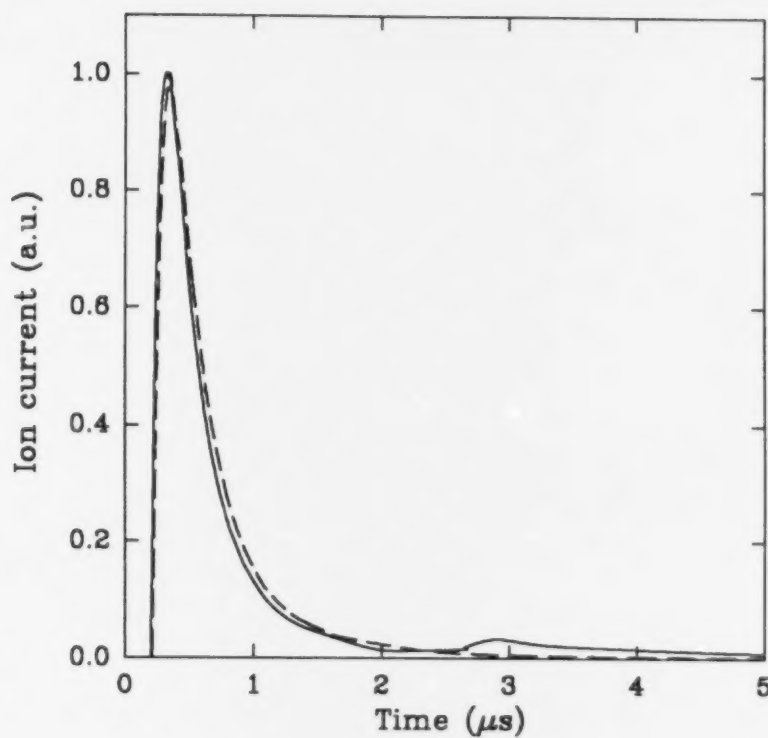


Figure 7.12 Simulated aluminum ion currents at 25 cm from the target surface by using plasma conditions calculated by Medusa with (solid line) and without (dashed line) radiation transport. The plasma is produced by a laser pulse of $\lambda = 0.25 \mu\text{m}$, $\tau = 1.5 \text{ ns}$ and $I_{\text{pk}} = 2 \times 10^{13} \text{ W cm}^{-2}$.

ion current with a peak ion velocity much faster than that of the measured ion current.

7.9 Summary of Charge State Distributions

The measured and calculated charge distributions have been compared. The calculated distributions were sensitive to the values of the three body recombination and the amount of heat returned to the plasma. However, using the values outlined in chapter 2 for the various rate coefficients and the starting plasma conditions from the Castor simulation reasonable agreement was found between the measured and observed distributions. However, discrepancies were observed in the late time charge states and the partial failure to predict the two peak behavior in aluminum. This discrepancy may indicate inaccuracies in recombination model and failure to include radiation energy transport in the modelling of the aluminum plasma.

CHAPTER EIGHT

CONCLUSION

Ion diagnostic measurements and numerical calculations from a hydro code that includes time dependent ionization and recombination physics are compared to understand the dominant physical processes involved in the expansion of a laser produced plasma to a large distance from the target surface. In order to accomplish this a comprehensive study of the hydrodynamics and energy balance of the UV laser produced plasma was carried out.

Experimental measurements were made of the expanding plasma produced by three different laser pulses: (1) a 2 ns KrF laser pulse at 248 nm at intensities of $10^{11} - 10^{13}$ W/cm², (2) a 20 ns KrF laser pulse at a laser intensity of 2×10^{11} W/cm² and (3) a 25 ns ruby laser pulse at laser intensities of 1×10^{11} to 2×10^{11} W/cm². In order to characterize the charge states of the plasma an electrodynamic multichannel analyzer was developed and employed for the study. The mass ablation rate, momentum balance and absorbed energy of the plasma were characterized using plasma calorimeters, Faraday cups and an integrating sphere. Plasma temperature measurements were made using continuum emission diagnostics.

At early time the plasma flow is planar which is then followed by a spherical expansion. After the laser is turned off and the plasma has expanded a few mm away from the target surface, most of the internal energy in the spherically expanding plasma is converted into kinetic energy of expansion. For low Z targets, most of the absorbed energy is converted into kinetic energy of ions at a large distance from the target surface which is confirmed by plasma energy measurements. The plasma under this condition can be described as undergoing an adiabatic spherical expansion.

Laser energy is found to be almost entirely absorbed by the targets ($> 90\%$). The plasma temperature and its intensity scaling deduced from x-ray continuum emission and ion diagnostics agree reasonably well with a steady state ablation model based on inverse bremsstrahlung absorption when the effect of energy invested in ionizing the plasma is included. The high level of laser absorption and the agreement of temperature scaling with the model are consistent with inverse bremsstrahlung being the dominant laser absorption mechanism. The sum of measured x-ray energy, kinetic energy and thermal energy of ions equals the absorbed laser energy indicating that the understanding of energy balance is reasonable. Reasonable agreement between the experimental angular distribution of plasma energy and that predicted by a plume model indicates the plume model is probably valid in describing the transition of planar to spherical flow for an expanding laser plasma.

Scaling laws for the mass ablation rates and ablation pressures from aluminum plasma based on ion diagnostics are obtained at $0.25 \mu\text{m}$ for intensities ranging from $10^{10} - 10^{13} \text{ W cm}^{-2}$. The results are in reasonable agreement with a 1 D steady state ablation model and results from other laboratories.

Energy transport measurement for gold plasma indicate the results are consistent with a two region plasma model consisting of a high temperature electron thermal transport region and a low temperature high density radiation transport region.

Recombination kinetics calculations indicate there is a time delay in ionizing the plasma early in the laser pulse. This is followed by a freezing of charge state at the end of the laser pulse. The calculated charge distributions were sensitive to the values of the three body recombination and the amount of heat returned to the plasma.

Reasonable agreement was found between the measured and

calculated distributions. However, discrepancies were observed in the late time charge states and partial failure to predict the two peak behavior in aluminum. This discrepancy may indicate inaccuracies in the recombination model and failure to include radiation energy transport in the modeling of the aluminum plasma.

Complete 2D simulations that includes radiation transport and time dependent ionization and recombination are probably required to resolve some of the remaining discrepancies between experimental observation and simulation of aluminum ion current.

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APPENDIX A

This appendix contains a partial program listing of the recombination kinetic code described in section 2.4.


```

C .....
C* Command for example this program
C* using Microsoft Fortran V4.0 or V5.0.
C*
C* fl G2.FPB87.A0x (G2.FPB87.A0x) for driver for output for
C* molibc for molibc for molibc for
C* .....
C* ID hydrodynamic code including ionization and recombination kinetics
C*
C* This code use MOL1D to integrate the 3 fluid equations and
C* the ionization and recombination rate equations
C*
C* The following units are used in the code:
C* [l] -> cm [m] -> ug [t] -> s
C* therefore, [rho] -> ug/cm3
C* [v] -> cm/s
C* [e,P] -> MerG/cm3
C*
C* The following conversion formula is used in the code
C* ni = Nav*rho/A [cm-3] rhoen = A*1e6/Nav [ug/cm3]
C* T = 18.77*(E-0.5*rho**2)/rho [eV]
C* E = 0.5*rho*v**2*2.4e-18*[cm-3]*[eV] [MerG/cm3]
C*
C* Brief introduction of subroutine MOL1D
C*
C* Method of lines integration package for systems of poles
C* UT = G(X,T,U,X,UX,FX), F = F(X,T,U,X,UX)
C*
C* The arguments of mol1d described how the system of n pole
C* partial differential equations are to be solved, the initial
C* conditions for the pole enter the routine in an array uz defining
C* the solution on a mesh xm of npts points, the user must describe
C* the kind of boundary conditions (kbc) for the equation and has
C* the option of describing the type of equation (keqp) that is
C* to be integrated (e.g. hyperbolic or parabolic, linear or
C* nonlinear), the mod approximation to be used is determined by the
C* temporal integration method (meth), the error per time step (epe)
C* and the spatial method order (mord), the user has the option of
C* allowing the package to choose a time integration method based on
C* the results of previous numerical testing, the equations are
C* integrated from the initial time tint to the final time tfat
C* with the output at user prescribed times in the array tout, the
C* character of the output (mout) may be either printed or graphics,
C* the printed output is written on tape 3 and the graphics use
C* plot88 or other graphic routines.
C* .....
C*
C* The mol1d package is composed of twenty fortran subroutines
C* including a modified version of the gear ode package.

```

```

C The subroutines in mol1d have the following names and perform
C the indicated tasks.
C MOL1D Driver routine for the package. It allocates storage,
C performs initialization tasks and monitors the progress
C of the time integration.
C CKINP Checks the validity of the input parameters
C SOLPRT Prints the solution at the user prescribed times
C SOLPLT Plots the solution at the user prescribed times
C SOLIMP Prints all relevant space derivatives of u and error
C estimates for the temporal integration, used for
C debugging and if a failure flag is returned from
C GEARB.
C DIFFUN Defines the time derivatives of u for integration
C package GEARB.
C BEVAL Evaluates the boundary conditions and incorporates
C them into the equation by cubic extrapolation to
C fictitious points outside the region of integration.
C UDER Evaluates the derivatives of u with respect to x.
C PDER Evaluates the derivatives of f with respect to x.
C FFTM42 A fast fourier transform routine
C FTFLT Filters high modes of the discrete fourier transform.
C COMPAT The user may call this routine to compute the
C accuracy in calculations when an exact solution is
C available.
C DRIBEB Driver routine for the gearb integration package.
C Other routines in GEARB INCLUDE INTERP, STIFFB, COSET,
C PSETB, DECB, SOLR, AND PDB.
C
C The temporal integration uses a. c. hindmarshs ode package
C GEARB solution of codes having banded jacobian
C u. of call: report ucid-30059, march 1975
C .....
C Input parameters
C
C The arguments of mol1d are defined as follows
C
C NPDE = Number of partial differential equations
C NPTS = Number of spatial mesh points
C KEQN = Descriptor of the kind of equation that is to be
C solved. This variable is used to pick the initial
C time step and the ode method if the default option
C is chosen. KEQN is greater than zero if the
C equation is linear and less than zero otherwise,
C it has the following values and meanings
C + - 1 hyperbolic equation
C + - 2 parabolic equation
C + - 3 equation in conservation form and/or solution
C has discontinuities
C - 4 nonstandard equation type

```

C KBC = Descriptor of the kind of boundary conditions for
 C the equation with the following values and meanings
 C 0 User supplies boundary conditions in subroutine
 C pde
 C 1 Periodic boundary conditions
 C 2 Boundary conditions of the form $a \cdot u + b \cdot u_x = c$
 C for $b \neq 0$ then a , b and c may be functions
 C of t and u . Otherwise a and c may be functions
 C of t only. Also included in this option are
 C boundary conditions defined by an ode,
 C free boundaries, and boundary conditions
 C prescribed.
 C 3 Boundary conditions as in kbc = 2 except that
 C some boundary conditions are defined by a pde
 C requiring spatial derivatives not used in the
 C main pde. All derivatives are evaluated at the
 C boundary in this option eliminating the need to set
 C MORD to a non zero value for those derivatives
 C needed only at the boundary.
 C
 C METH = The ode method flag used to set the calling arguments
 C in GEARB. If either of the default options is used,
 C METH returns equal to the method selected.
 C This method is chosen by the values of KEQN, KBC,
 C NPTS and NPDE. METH has two decimal digits.
 C MODE and MJAC (METH = 10*MODE + MJAC). MODE is
 C the ode method indicator with the following values
 C and meanings
 C
 C MODE = 0 Mol1d selects one of the following methods
 C MODE = 1 Adams-Bashford-Moulton methods
 C MODE = 2 Gear stiffly stable backward difference methods
 C
 C MJAC is the corrector iteration method
 C indicator, with the following values and meanings
 C
 C MJAC = 0 MOL1D selects one of the following methods
 C MJAC = 1 Jacobian approximated by the identity matrix
 C (function iteration)
 C MJAC = 2 Jacobian approximated by a diagonal matrix
 C MJAC = 3 Jacobian approximated by a banded matrix, band
 C width = NPDE (for stiff equations)
 C MJAC = 4 Jacobian approximated by a banded matrix,
 C band width = 3*NPDE (for stiff equations)
 C MJAC = 5 Jacobian approximated by a banded matrix,
 C band width = 5*NPDE (for stiff equations)
 C
 C EPS = Time integrator local error tolerance parameter
 C estimates of the time step error at the spatial
 C grid points for all the pdes is kept less than eps
 C in the root-mean-square (rms) norm. EPS is also
 C
 C used to choose the initial step.
 C
 C MOUT = An input parameter to indicate the character of the
 C output. If MOUT is less than zero then the solution
 C is printed at each time step for IABS (MOUT) steps.
 C TOUT(I) returns the time of the i-th time step.
 C If MOUT is greater than equal to zero, it has
 C the following values and meanings
 C 0 No output is printed or plotted
 C 1 The solution is printed at the preselected times
 C in the array TOUT. The printing is in subroutine
 C SOLPRT on file 3. This subroutine may be
 C rewritten by the user to change the output format.
 C 2 The solution and its time derivative are printed
 C at the times in TOUT.
 C 3 The solution is printed and plotted at the times
 C in TOUT. The plotting is in subroutine solpl.
 C an alternative subroutine may be supplied by the
 C user if one wishes to provide his/her own graphics
 C routine.
 C 4 The time derivative of the solution is printed
 C and the solution is both printed and plotted at
 C the times in tout.
 C 5 The solution is plotted at the times in TOUT.
 C
 C TOUT = An input array of the times that the solution is
 C to be printed or plotted. If MOUT .gt. 0, the
 C solution is interpolated at these times from its
 C values at the actual integration times by Taylor
 C series. If MOUT .lt. 0 then the first IABS (MOUT)
 C printout times are returned in tout.
 C
 C MORD = A matrix dimensioned mord (NPDE,3) indicating which
 C derivatives of the pde variables are needed and
 C what type of approximation is to be used.
 C The derivative calculations done by the package all
 C assume an equally spaced mesh, if an unequally spaced
 C mesh is used the set mord(i,j)=0 or a terminal error
 C message will result.
 C MORD(1,1), MORD(1,2) and MORD(1,3) refer to
 C UX(I), UXX(I) and FX(I) respectively. If the
 C derivatives are needed only at the boundary set KBC=3
 C rather than setting mord to a non zero value. MORD has
 C the following possible values and meanings
 C 6 Sixth order centered differences
 C 4 Fourth order centered differences
 C 2 Second order centered differences
 C 1 First Fourier transform approximation
 C 0 Derivative is not to be calculated by the package.
 C It may still be calculated by the user in
 C subroutine PDE. Note that no derivatives can

```

C - be calculated by the package on an unequally spaced mesh.
C -1 Fast Fourier transform with linear filtering of the higher modes.
C -2 Second order differences oriented into the characteristics of u = fx. The centered differences are used if the characteristics are perpendicular to the x axis.
C -4 Fourth order centered differences or unsymmetric third order differences oriented into the characteristics of u = fx.
C
C TINT = The initial value of the independent variable t
C
C TLAST = The final integration time when control is returned to the calling program. TLAST must be greater than or equal to TINT. If the ode integrator fails to converge during integration TLAST returns equal to the last integration time.
C
C UZ = A matrix dimensioned UZ(NPDE, NPPTS) of the dependent variable U. On input UZ(LJ) contains the value of i-th pde component at the j-th mesh point at time T = TINT. ON output, UZ returns the value of the solution at time = TLAST.
C
C XM = AN array of length nppts containing the mesh points. The mesh points must be evenly spaced. If the derivatives are calculated by the package, XM need not be evenly spaced if MORD(LJ)=0 for all i,j. The derivatives can then be computed in subroutine PDE by the user
C
C KMOL = An indicator flag used for both input and output. On input it has the following values and meanings k.h.0 Calls a debugging dump on the first and last time step. In debugging dump the solution, all relevant spatial and temporal derivatives and error estimates are printed. K = 0 No dumps are called unless an error is encountered
C
C On output, KMOL returns 0 If the integration was successful K.L.T. 0 The integration was halted after encountering a fatal error. A message is printed and K returns with the following values and meanings
C -1 Integration was halted after failure to pass the error test even after the time step, dt, was reduced by a factor of 10**10 from its original value.
C

```

```

      IMPLICIT REAL*8(A-H,O-Z),INTEGER*4(I-N)
      CHARACTER COM*30,BCFILE*12
      PARAMETER (NEQ=17,NIX=2550,NGRID=43350)
      DIMENSION UZ(NEQ,NDX),XM(NDX,MORD1,NEQ,5),TOUTI(201)

C...
      COMMON/ARFVTS/ARFV(NEQ)
      COMMON/GRIDS/NPTS
      COMMON/TIMEBC/NIBC,TBC(50) UBC(NEQ,50)
      COMMON/GMETRY/NALPHA
      COMMON/MINU/UMIN(NEQ)
      COMMON/IC/U(NEQ)
      COMMON/RATES/NINZ
      COMMON/ATOMIC/NZ,AMASS,XI(14)

C... open a outputfile
      OPEN(2,FILE='khyd.out')
      OPEN(3,FILE='khyd.dmp')
      OPEN(4,FILE='khyd.par',STATUS='OLD')

C... read input parameters
      READ(4,*)
      READ(4,803)COM
      READ(4,*)
      READ(4,803)BCFILE
      READ(4,*)
      READ(4,*)EPS,KBC,KEON,KMOL,METH,MOUT
      READ(4,*)
      READ(4,*)NPDE,NPTS,X0,XEND
      READ(4,*)
      READ(4,*)NALPHA
      READ(4,*)
      READ(4,*)ARFV(I),I=1,NPDE)
      READ(4,*)
      READ(4,*)NINZ
      READ(4,*)
      READ(4,*)NSTEP,T0,TEND
      READ(4,*)
      READ(4,*)U(0,I),I=1,NPDE)
      READ(4,*)
      READ(4,*)UMIN(I),I=1,NPDE)
      READ(4,*)
      READ(4,*)XNDROP
      CLOSE(4)

C...
      DO 10 I=1,NPDE
        MORD1(I)=0
        MORD1(2)=2
        MORD1(3)=4
        MORD1(4)=4
      10 MORD1(5)=4
      C... ionization potentials
      NZ=NPDE-3
      IF(NZ.EQ.7) THEN
        AMASS=12
        XI(1)=10.6
        XI(2)=26.05
      ELSE
        AMASS=27.
        XI(1)=5.714
        XI(2)=18.66
        XI(3)=29.15
        XI(4)=116.2
        XI(5)=161.5
        XI(6)=206.7
        XI(7)=252.0
        XI(8)=297.9
        XI(9)=343.4
        XI(10)=427.1
        XI(11)=471.1
        XI(12)=1984.
        XI(13)=2211.
      ELSE
        WRITE(*,*)'Error! Element other than Al or C'
        stop
      ENDIF

C... define output timesteps
      DT=(TEND-T0)/DBLE(NSTEP-1)
      DO 20 I=1,NSTEP
        TOUT(I+1)=T0+DT*DBLE(I-1)
      20 CONTINUE
      NTOUT=NSTEP+1
      TOUT(1)=0.

      TINT=TOUT(1)
      TLAST=TOUT(NTOUT)
      C... define mesh and define initial conditions
      XL=XEND-X0
      DX=XL/DBLE(NPTS-1)
      DO 30 I=1,NPTS
        XM(I)=X0+DBLE(I-1)*DX
      30 UZ(K,I)=U0(K)*(XDROP/XM(I))**3
      ELSE
        DO 40 K=1,NPDE
          UZ(K,I)=U0(K)
        40 UZ(K,I)=U0(K)
      ENDIF
      C... do not allow it less than a preset value
      DO 45 K=1,NPDE
        IF(UZ(K,I).LT.UMIN(K)) UZ(K,I)=UMIN(K)
      45 IF(UZ(K,I).LT.UMIN(K)) UZ(K,I)=UMIN(K)
      50 CONTINUE
      C... get time dependent boundary conditions
      OPEN(5,FILE=BCFILE,STATUS='OLD')
      WRITE(*,*)'Reading b.c. ....'

```



```

      UT(2,J)=FX(2,J)/(XM(J)*NA)
      & +AFV(2)*(UXX(2,J)+DBLE(NA)*UX(2,J)*XM(J))
      & +DBLE(NA)*P/XM(J)
      UT(3,J)=FX(3,J)/(XM(J)*NA)
      & +AFV(3)*(UXX(3,J)+DBLE(NA)*UX(1,J)*XM(J))+ERTRN
    ENDIF
  C
  I=1
  UT(1+3,J)=V*FX(1+3,J)
  & +XNE*(ALPHA(1+1)*F(1+1)-S(1)*F(1))
  & +AFV(1+3)*UXX(1+3,J)
  DO 70 I=2,NZ-1
    UT(1+3,J)=V*FX(1+3,J)
    & +XNE*(S(1)*F(1)-ALPHA(1+1)*F(1+1))
    & +S(1+1)*F(1+1)+ALPHA(1+1)*F(1+1)
    & +AFV(1+3)*UXX(1+3,J)
  70 CONTINUE
  I=NZ
  UT(1+3,J)=V*FX(1+3,J)
  & +XNE*(S(1)*F(1)-ALPHA(1)*F(1))
  & +AFV(1+3)*UXX(1+3,J)
200 CONTINUE
  C
  IX=NPTS
  RETURN
  END
C+++end pde
C .....
C+++start func
SUBROUTINE BNDRY(T,U,L,AL,BL,CL,UR,AR,BR,CR,NPDE)
  IMPLICIT REAL*8(A-H,O-Z),INTEGER*4(I-N)
  PARAMETER (NEQ=17,NIX=250,NGRID=43350)
  DIMENSION UR(NPDE),AR(NPDE),BR(NPDE),CR(NPDE)
  DIMENSION UL(NPDE),AL(NPDE),BL(NPDE),CL(NPDE)
  DIMENSION UNBC(50)
  COMMON/TIMEBC/TIBC,TBC(50),UBC(NEQ,50)
  COMMON/MINU/UMIN(NEQ)
  COMMON/IC/UQ(NEQ)
  C .....
  C.....
  DO 5 J=1,NPDE
    5 AL(J)=1.0
    DO 20 J=1,NPDE
      DO 10 K=1,NIBC
        10 UNBC(K)=UBC(J,K)
      CALL NCSPLN(NIBC,TBC,UNBC,T,CL(J))
    20 CONTINUE
    C..... do not allow h.c. less than a preset value
    DO 50 J=1,NPDE
      50 IF(CL(J).LT.UMIN(J))CL(J)=UMIN(J)
    C..... free h.c.
    DO 60 J=1,NPDE
      60 BR(J)=1.0
  C .....
  RETURN
  END
C+++end body
C .....
C+++start acapln
C/* Subroutine to do cubic spline
SUBROUTINE NCSPLN(NPTS,XA,YA,X,Y)
  IMPLICIT REAL*8(A-H,O-Z),INTEGER*4(I-N)
  DIMENSION XA(NPTS),YA(NPTS)
  DIMENSION A(100),B(100),C(100),D(100)
  DIMENSION CL(100),CM(100),CZ(100)
  DO 10 I=2,NPTS-1
    A(I)=YA(I+1)*XA(I)-XA(I-1))
    A(I)=A(I)-YA(I)*XA(I+1)-XA(I-1))
    A(I)=A(I)+YA(I-1)*XA(I+1)-XA(I)
    A(I)=A(I)+3*((XA(I+1)-XA(I))*XA(I)-XA(I-1)))
  10 CONTINUE
  C CL(1)=1.

```

```

C.....
C++++start ionzr2
C.....
C..... Subroutine calculate the ionization rate from:
C..... Colomant (1973).
C.....
SUBROUTINE IONZR2(N,XI,TEMP,S)
IMPLICIT REAL*8(A-H,O-Z),INTEGER*4(I-N)
DIMENSION S(14),XI(14)
C
DO 10 I=1,N-1
CALL SIGMA(N-1,SIGMA)
S(I)=9.D-6*SIGMA*TEMP**0.5/XI(I)**2*(4.88*TEMP/XI(I))
S(I)=S(I)*DEXP(-XI(I)/TEMP)
10 CONTINUE
S(N)=0.
RETURN
END
C++++end ionzr2
C.....
C++++start sigma
C.....
C..... Subroutine to calculate the number of electrons in the outer
C..... shell for ionization state z (1 for hydrogenic like ions)
SUBROUTINE SIGMA(NZ,I,SIGMA)
IMPLICIT REAL*8(A-H,O-Z),INTEGER*4(I,N)
NDIFF=NZ-1
N=1
SIGMA=1.0
IF(NDIFF.GT.2)N=2
IF(NDIFF.GT.10)N=3
IF(NDIFF.GT.28)N=4
IF(NDIFF.GT.36)N=5
IF(NDIFF.GT.54)N=6
SIGMA=NZ-2*(N-1)**2.1
RETURN
END
C++++end sigma
C.....
C++++start idcr
C.....
C..... Subroutine calculate the radiative recombination rate.
C..... (D.L. Book, NRL Plasma Formulary (1982).
C.....
SUBROUTINE RDCR(N,XI,TEMP,ALPHAR)
IMPLICIT REAL*8(A-H,O-Z),INTEGER*4(I,N)
DIMENSION ALPHAR(14),XI(14)
C
ALPHAR(1)=0.
DO 10 I=2,N
FI=DBLE(I-1)
ALPHAR(I)=DBLE(43)+DBLE(.5)*DLOG(XI(I-1)/TEMP)
ALPHAR(I)=ALPHAR(I)+DBLE(.47)*(XI(I-1)/TEMP)**(1./3.)

```

```

C.....
C.....
CM(1)=0.
CZ(1)=0.
C
DO 20 I=2,NPTS-1
CL(I)=2*(XA(I-1)-XA(I-1))*(XA(I)-XA(I-1))*CM(I-1)
CM(I)=(XA(I-1)-XA(I))/CL(I)
CZ(I)=(A(I)-XA(I)-XA(I-1))*CZ(I-1)/CL(I)
20 CONTINUE
CL(NPTS)=1.
CZ(NPTS)=0.
C(NPTS)=0.
C
DO 30 J=2,NPTS-1
C(J)=CZ(J)-CM(J)*C(I+1)
BI(J)=(YA(J-1)-YA(J))/(XA(J-1)-XA(J))
BJ(J)=BI(J)-XA(J-1)-XA(J)*C(I+1)+2.*C(I)/3.
DI(J)=(C(I+1)-C(J))/(3*(XA(J-1)-XA(J)))
30 CONTINUE
DO 40 J=2,NPTS
IF(XA(J).GT.X)THEN
Y=YA(J-1)+BJ(J)*(X-XA(J-1))+C(J-1)*(X-XA(J-1))**2
Y=Y+DI(J)*(X-XA(J-1))**3
GO TO 100
ENDIF
40 CONTINUE
I=NPTS
Y=YA(J-1)+BJ(J)*(X-XA(J-1))+C(J-1)*(X-XA(J-1))**2
Y=Y+DI(J-1)*(X-XA(J-1))**3
100 RETURN
END
C++++end ncopl
C.....
C++++start ionzr1
C.....
C..... Subroutine calculate the ionization rate. An analytic
C..... expression which was derived by averaging the ground state
C..... ion over a Maxwellian electron distribution (D.L. Book,
C..... NRL Plasma Formulary (1982)).
C.....
SUBROUTINE IONZR1(N,XI,TEMP,S)
IMPLICIT REAL*8(A-H,O-Z),INTEGER*4(I,N)
DIMENSION S(14),XI(14)
C
DO 10 I=1,N-1
S(I)=DSQRT(TEMP/XI(I))
S(I)=S(I)/DSQRT(XI(I)**3)*(DBLE(6)+TEMP/XI(I))
S(I)=S(I)*DEXP(-XI(I)/TEMP)*1.0D-5
10 CONTINUE
S(N)=0.
RETURN
END
C++++end ionzr1

```



```

      ALPHAR(I)=ALPHAR(I)*F*DSQRT(XI(I-1)/TEMP)*S.2D-14
10 CONTINUE
      RETURN
      END
C+++end rdxr
C.....
C+++start b3cr
C.....
C++++ Subroutine calculate the three body recombination rate
C++++ using the analytic expression by Veselovski
C++++ (I.S. Veselovski, Sov. Phys. Tech. Phys. 14, 193 (1969)).
C.....
SUBROUTINE B3RCR(N,TEMP,ALPHA3)
      IMPLICIT REAL*8(A-H,O-Z),INTEGER*4(I-N)
      DIMENSION ALPHA3(14)
      ALPHA3(1)=0.
      DO 10 I=2,N
         FI=DBLE(I-1)
         ALPHA3(I)=(FI**3)*DLOG(DSQRT(FI**2+DBLE(1))) / DSQRT(TEMP**9)
         ALPHA3(I)=ALPHA3(I)*8.75D-27
10 CONTINUE
      RETURN
      END
C+++end b3cr
C.....
C+++start rheat
C.....
C++++ Subroutine calculate the energy return to the plasma
C++++ because of recombination. The expression used is derived
C++++ by Zel'dovich and Raizer (Ya.B. Zel'dovich and
C++++ Yu.P. Raizer, Physics of Shock Waves and High Temperature
C++++ Hydrodynamic Phenomena, (1966)
C.....
SUBROUTINE RCHEAT(N,TEMP,ZAVE,DNION,ERC)
      IMPLICIT REAL*8(A-H,O-Z),INTEGER*4(I-N)
      DIMENSION ERC(14)
      ERC(1)=0.
      DO 10 I=2,N
         FI=DBLE(I-1)
         ERC(I)=(ZAVE**(1.6)) / (FI**(2.3))
         ERC(I)=ERC(I)*(ZAVE*DNION)**(1.6))
         ERC(I)=ERC(I)*(TEMP**(1.12))**1.1D-3
10 CONTINUE
      RETURN
      END
C+++end rheat
C.....
C+++start avegz
C.....
C++++ Subroutine calculate the average Z
C.....
SUBROUTINE AVEGZ(N,F,ZAVE)

```

```

      IMPLICIT REAL*8(A-H,O-Z),INTEGER*4(I-N)
      DIMENSION F(14)
      ZAVE=0.0
      DO 20 I=1,N
         ZAVE=ZAVE+F(I)*DBLE(I-1)
20 CONTINUE
      RETURN
      END
C+++end avegz
C&&&&
      .

```